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THE XY MODEL FOR PHASE TRANSITIONS -
SOME STATIC AND DYNAMIC RESULTS

BY



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A THESIS

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The undersigned certify that they have read, and recommend to the Faculty of Graduate Studies for acceptance, a thesis entitled "The XY Model for Phase Transitions - Some Static and Dynamic Results", submitted by Ruth Viterbo Ditzian in partial fulfillment of the requirements for the degree of Doctor of Philosophy.

Date

July 6, 1970..

ABSTRACT

Exact high temperature series expansions have been obtained for various thermodynamic quantities for the spin one half XY model of ferromagnetism or of a quantum lattice fluid.

The static quantities calculated include the initial parallel susceptibility and the fourth order fluctuation both for the face centered cubic and for the triangular lattice. Analysis of the series for these two quantities by standard ratio and Pade approximant methods yields estimates for the critical temperature variable, $K_c = J/kT_c$, and the critical indices γ and γ_2 respectively. For the face centered cubic lattice $K_c = 0.221 \pm 0.001$, $\gamma = 1.32 \pm 0.04$ and $\gamma_2 = 4.64 \pm 0.10$. For the triangular lattice $K_c = 0.67 \pm 0.01$, $\gamma = 1.50 \pm 0.02$ and $\gamma_2 = 5.4 \pm 0.5$. These results coupled with scaling theory yield all critical indices for the spin one half XY model. On the triangular lattice the results obtain if we assume a phase transition but do not constitute conclusive evidence of the existence of such a transition. The results for the face centered cubic lattice compare very well with experimental results for the λ transition in liquid helium.

The dynamical calculations have been done for two quantities, the autocorrelation and the frequency

dependent parallel susceptibility, both on the face centered cubic lattice. The autocorrelation can be expanded in a power series in time. We have calculated the first, second, third, fourth, sixth and seventh coefficients as power series in T^{-1} . We find that the first, second and third coefficients diverge at K_c with exponents $\gamma_1^A = 0.10 \pm 0.01$, $\gamma_2^A = 0.24 \pm 0.05$ and $\gamma_3^A = 0.35 \pm 0.05$ respectively.

The frequency dependent susceptibility can be expanded in a power series in the reciprocal of the frequency, ω^{-1} . Coefficients of odd powers vanish; coefficients of even powers have been obtained as power series in the reciprocal temperature variable, K . We find that the second moment diverges at K_c with exponent $\gamma_1 = 0.10 \pm 0.01$ and the fourth moment vanishes with exponent $\gamma_3 = -0.86 \pm 0.05$. Assuming that a relation $\chi(\omega, T) = [(T_c - T)/T_c]^{-\gamma} f(\omega/[(T - T_c)/T_c]^{\Delta_s})$ holds near T_c for all ω we find that $\Delta_s = 0.58 \pm 0.10$. It is important to observe that $\Delta_s < \gamma$.

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CHAPTER I

INTRODUCTION

This thesis describes a study of the thermodynamical properties above the critical temperature for the spin $\frac{1}{2}$ XY model. The method of exact high temperature series expansion is used. The XY model is quantum mechanical and yet calculations with it are relatively easy. The spontaneous magnetization, which is the order parameter, lies in the xy plane. Therefore the order parameter in the XY model does not commute with the Hamiltonian and can exhibit dynamical behaviour. This is not the case in the Ising and Heisenberg models, where the order parameter is constant in time and changes in the Hamiltonian have to be introduced to allow dynamical behaviour of the magnetization. Accordingly this thesis is divided into two parts. The first deals with static critical phenomena and the second treats dynamical critical phenomena.

In Section 2.2 the XY model is introduced. We shall briefly summarize here the reasons for the interest in this model. In 1956 Matsuda and Matsubara proposed the anisotropic Heisenberg model as a possible model for an interacting Bose fluid. The XY model represents a Bose fluid with hard core repulsive interaction only.

The classical anisotropic Heisenberg model has been studied by Jasnow and Wortis (1968) for various values of the anisotropy parameter. They concluded that the critical behaviour for the classical model depends only on the dimensionality of the lattice and the symmetry of the ground state. That is, the critical indices are constant over ranges of the anisotropy parameter and change discontinuously when the symmetry of the ground state changes. The XY behaviour persisted as long as the ground state was invariant under rotation in the plane. This is true when the XY part of the interaction is stronger than the Z or Ising part.

Jasnow and Wortis studied only the spin infinity, or classical, case but the conjectured behaviour is supported for the quantum mechanical case by the work of Barouch and McCoy (1969). Barouch and McCoy found, for the anisotropic XY model in one dimension, that the nature of the decay of the spin correlations changed when changes in the anisotropy parameter and in the applied field affected the symmetry of the ground state. It seems therefore reasonable to expect that the critical indices of the spin $\frac{1}{2}$ anisotropic Heisenberg model will be equal to the indices of the XY model as long as the

XY interaction is dominant. The indices we obtain are expected then to fit the λ transition of liquid helium, and some insulating magnetic crystals for which $J_{\perp} \gg J_{\parallel}$, that is the transverse interaction is much stronger than the parallel.

The quantities usually expanded in high temperature series are free energy, the susceptibility and the even derivatives of the free energy with respect to the field at zero field. Betts, Elliott and Lee (1970) calculated the free energy and fluctuation series for a variety of lattices. Because the magnetization and the Hamiltonian do not commute, the fluctuation, $\frac{1}{N} \langle M_x^2 \rangle$, is not equal to the susceptibility. Falk and Bruch (1969) proved however that the critical index of the fluctuation is equal to γ , the index of the susceptibility. The proof involved only rigorous thermodynamical inequalities. The equality is only asymptotically true close to the critical temperature. It is therefore of interest to calculate the susceptibility to see numerically how the coefficients approach those of the fluctuation series.

From the fluctuation series on the triangular lattice no conclusions could be drawn as to the existence of a

transition and it was hoped that the static susceptibility would behave better.

The other static quantity calculated was the fourth order parallel fluctuation

$$Y_2 = \frac{1}{N} \left\{ \frac{3}{2} \langle M_x^2 \rangle^2 - \frac{1}{2} \langle M_x^4 \rangle \right\} .$$

By similar arguments to those used for the fluctuation by Falk and Bruch (1969), Y_2 diverges as the fourth derivative of the free energy with respect to the field in the x direction, thereby giving us an estimate for the gap parameter. Now two critical exponents and scaling yield estimates for all critical indices.

Before leaving the static quantities we should mention that the limit at zero frequency of the dynamical susceptibility is the quasistatic (called also adiabatic after Kubo) susceptibility while ours is the isothermal. They are equal when the system is ergodic. This question arose in the dynamical calculations for the Ising model of Allan and Betts (1968) and Essam and Garelick (1968). We shall discuss ergodicity in Chapter V but mention here the conclusion that in the XY 3-dimensional case we believe the two are equal.

Dynamic critical phenomena are at a much earlier stage of theoretical investigation than static critical phenomena.

Scaling laws have less success with static correlations than they had with the thermodynamical quantities. For the dynamical correlations the assumptions used for the static correlations are kept and generalized, therefore scaling is on a weaker basis.

There is no exactly soluble model with physical critical behaviour that can be used to check dynamical theories the way the 2d Ising model is used in static theory. The only soluble model is the one dimensional XY model. Niemeijer (1967), Barouch, McCoy and Dresden (1969) and Suzuki (1969) have calculated the time dependent correlations and magnetization of the one dimensional XY model. This model has the disadvantage of being non ergodic. The limit $t \rightarrow \infty$ yields a magnetization which is not the equilibrium magnetization in the z direction. The magnetization in the x direction can be expected to be ergodic and of interest but cannot be calculated exactly with present methods.

The kinetic Ising model in 2 dimensions was investigated by Suzuki (1969) using high temperature series expansions. Suzuki found that the critical index of slowing down is different from the index of the static

susceptibility. Dynamical molecular field theory (Van Hove (1959), Suzuki and Kubo (1968)) predicts equality for the two indices.

Recently McFadden and Tahir Kheli (1970) calculated the zeroth, second and fourth moments of the spectral function as functions of the momentum transfer. Only the leading term in T^{-1} was obtained.

We have calculated the frequency dependent susceptibility $\chi(\underline{Q}, \omega)$. We have obtained a double series in powers of T^{-1} and ω^{-1} . Analysis of the series is hampered by the fact that coefficients of powers of T^{-1} are infinite series in powers of ω^{-1} and vice versa. We find that the second moment of the susceptibility diverges at $T > T_c$. For the kinetic Ising model and the isotropic Heisenberg model all moments are finite, as was rigorously proved by Suzuki (1969) and Mermin and Wagner (1966) respectively. Though the proofs could not be generalized to the XY model the divergence was not expected. Assuming scaling we estimate the critical index of slowing down, and we find that in this case it is speeding up. We calculated the auto-correlation and with a scaling assumption estimate its relaxation time.

CHAPTER II

STATIC CRITICAL PHENOMENA AND THE XY MODEL

2.1 Critical phenomena and critical exponents

Phase transitions have been studied by physicists since the day man discovered that water boiled and froze.

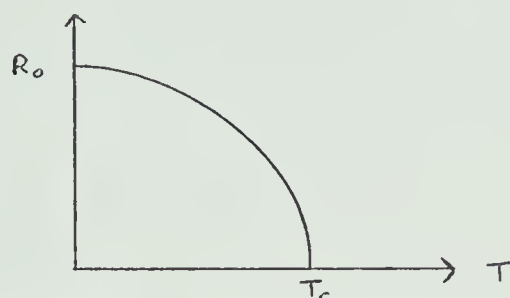
A great variety of such phenomena, while of completely different origin, have a very similar quantitative behaviour and can be studied using the same mathematical models.

Some examples of critical phenomena are those which occur in monomolecular fluids near the critical point, in liquid helium near the λ temperature, in binary alloys near the order-disorder transition and in ferromagnets in the vicinity of the Curie temperature.

In all these systems by varying a thermodynamical parameter we reach a point where there is no phase transition, only one homogeneous phase exists and changes such as in density or magnetization are continuous.

For all those phenomena a long range order parameter R can be defined. At $T=0$ there is absolute order in the system. All spins are lined up parallel to each other in a ferromagnet, all He^4 atoms are in the ground superfluid state etc., so $R=R_0$; at infinite temperature $R=0$ as complete disorder must occur. In the systems with which we are concerned however the long range order disappears suddenly at and above some temperature defined as the critical point,

whereas for systems which do not exhibit phase transitions R decays to zero smoothly and only at $T=\infty$; for the transitions with which we deal the behaviour of R is



For example, the order parameter for a ferromagnet is the magnetization, and for superfluid it is the density of the superfluid fraction.

In the vicinity of this critical point there are large fluctuations of all thermodynamic quantities which can become macroscopic and the relevant derivatives of the thermodynamic potential, such as specific heat and susceptibility, diverge. Such macroscopic fluctuations cannot probe deeply into the microscopic details of the forces involved and this is the reason of the quantitative similarity among such diverse phenomena.

The determination of the asymptotic laws which describe the approach to the critical point has been the main problem in the study of critical phenomena. It may be that only the dimensionality and symmetry of the system are needed to predict the asymptotic laws of the appropriate variable. This is a unifying pleasing hypothesis (Jasnow, Wortis 1968; Kadanoff 1970).

(a) Magnetic systems

Ferromagnets are defined as those crystals which have a spontaneous magnetization below the Curie temperature T_c , that is which possess a magnetic moment even in the absence of an applied magnetic field H .

$$M_o(T) = \lim_{H \rightarrow 0+} M(H, T) \quad (2.1.1)$$

Above T_c is the paramagnetic region and the magnetization varies continuously with H through the line $H=0$.

The magnetic exponent β is defined by

$$M_o(T) \sim (T - T_c)^\beta \quad (2.1.2)$$

as T approaches T_c . This behaviour is well verified by the experimental data of Heller and Benedek (1965) and Senturia and Benedek (1966) on EuS and CrBr_3 .

The isothermal zero field susceptibility

$$\chi_T = \left(\frac{\partial M}{\partial H} \right)_{H=0} \quad (2.1.3)$$

diverges as the temperature approaches T_c

$$\chi_T \sim \begin{cases} (T - T_c)^{-\gamma} & T > T_c \\ (T - T_c)^{-\gamma'} & T < T_c \end{cases} \quad (2.1.4)$$

This relation defines the critical exponents γ and γ' and again experiments bear out this assumed behaviour for a

variety of materials: Ni, Fe, Gd, YtFeO₃ (see Miedema et al (1963)). The critical index δ is defined on the critical magnetic isotherm

$$H \sim |m|^\delta \quad T = T_c \quad . \quad (2.1.5)$$

The specific heat critical exponents α and α' are defined by

$$C_{H=0} \sim \begin{cases} (T-T_c)^{-\alpha'} & T < T_c \\ (T-T_c)^{-\alpha} & T > T_c \end{cases} \quad (2.1.6)$$

$\alpha=0$ is by definition a logarithmic singularity, for example EuO (Teaney 1966).

One more series of exponents Δ'_n can be defined (for $T < T_c$) by the relation:

$$\left(\frac{\partial^n F}{\partial H^n} \right)_T \sim (T_c - T)^{-\Delta'_n} \left(\frac{\partial^{n-1} F}{\partial H^{n-1}} \right)_T \quad (2.1.7)$$

where F is the free energy, and similarly Δ_n is defined for $T > T_c$.

For $n=1$, (2.1.7) does not define new exponents as they are related to α, β, γ' by

$$\Delta'_1 = 2 - \alpha' - \beta$$

$$\Delta'_2 = \beta + \gamma' \quad .$$

Above T_c odd derivatives vanish for $H=0$ because $M(-H, T) = -M(H, T)$.

Δ_n and Δ'_n cannot be measured experimentally for $n > 2$.

The exponents ν, ν' and η refer to the behaviour of the pair correlation function $\Gamma(r)$ and correlation length ξ observed in neutron scattering, first explained by Van Hove (1945)

$$\xi \sim \begin{cases} (T-T_c)^{-\nu'} & T < T_c \\ (T-T_c)^{-\nu} & T > T_c \end{cases} \quad (2.1.8)$$

$$\Gamma_c(r) \approx \frac{1}{r^{d-2+\eta}} \quad \begin{matrix} T = T_c \\ H = 0 \end{matrix} \quad (2.1.9)$$

where d is the dimensionality of the system.

The spin correlation function is defined by

$$\Gamma_{\alpha\beta}(\underline{r}, H, T) = \frac{\langle S_o^\alpha S_{\underline{r}}^\beta \rangle - \langle S_o^\alpha \rangle \langle S_o^\beta \rangle}{\frac{1}{3} S(S+1)} \quad \alpha, \beta = x, y, z \quad (2.1.10)$$

and Γ is the correlation function in the preferred direction at the critical point. The quasi elastic scattering intensity for a momentum transfer \underline{q} is:

$$\frac{I(\underline{q})}{I^0(\underline{q})} = 1 + \sum e^{i\underline{q} \cdot \underline{r}} \Gamma(\underline{r}) \quad (2.1.11)$$

The coherent scattering at zero field below T_c is

$$I_{\text{coh}}(0, T) \sim M_o^2(T) \sim (T-T_c)^{2\beta} \quad (2.1.12)$$

which gives another experimental way of measuring β .

The zero angle scattering is proportional to the susceptibility,

$$I(0) \sim \chi_T \sim (T-T_c)^{-\gamma} \quad H = 0 \quad (2.1.13)$$

$$T > T_c$$

Experiments were performed by Passel et al (1965), Balley et al (1967) on iron to measure γ .

From (2.1.9) and (2.1.11) we see

$$I_{\text{critical}}(\underline{k}) \sim \frac{1}{k^{2-\eta}} \quad \text{as } k \rightarrow 0 \quad . \quad (2.1.14)$$

Jacrot et al (1962), Passel et al (1965) measured η this way.

The spin correlation function decays as

$$\Gamma(\underline{r}, T) \sim \frac{e^{-\kappa r}}{r} \quad . \quad (2.1.15)$$

The correlation length is defined by

$$\xi = \kappa^{-1} \quad . \quad (2.1.16)$$

$\kappa_1(T)$ measures the slope of $I^{-1}(\underline{k}, T)$ against k^2 as $k \rightarrow 0$

$$\kappa_1^{-2} \propto \frac{\int r^2 \Gamma(\underline{r}) d\underline{r}}{1 + v^{-1} \int \Gamma(\underline{r}) d\underline{r}} \quad . \quad (2.1.17)$$

Near T_c we expect $\kappa_1 \sim \kappa$ if near the critical point only one divergent temperature dependent correlation length exists. The quantity v that is measured is the one pertaining to κ_1 .

This was a short summary of the definitions of critical indices for uniaxial ferromagnets; our interest is really in ferromagnets with a plane of easy magnetization but applying a small field in a direction lying in the plane enables us to treat that direction as the preferred axis of magnetization and all definitions still apply.

(b) Superfluid helium

At first sight the superfluid transition seems rather different from the transition in a uniaxial ferromagnet because of the existence of the λ line (see Figure 2.2). The confusion arises due to regarding the pressure on the superfluid as the analogue of the magnetic field and the density as the analogue of the magnetization. The correct analogy is between the superfluid order parameter (the square root of the superfluid density) and the magnetic order parameter (the axial magnetization). The generalized force conjugate to the superfluid order parameter and analogous to the axial magnetic field is unfortunately not physically realizable. The analogues of the density and pressure are more correctly in the uniaxial magnet, the perpendicular magnetization and the perpendicular magnetic field. We shall expect therefore

$$n_0(t) \sim (T - T_\lambda)^{2\beta} \quad . \quad (2.1.18)$$

n_0 is defined by

$$n_0 = |\psi|^2 = \lim_{(\underline{r} - \underline{r}') \rightarrow \infty} \langle \psi^\dagger(\underline{r}) \psi(\underline{r}') \rangle \quad (2.1.19)$$

$\Psi(\underline{r})$ being the wave function in second quantization formulation. Most simple theories of superfluidity assert

$$\rho_s \sim n_0$$

ρ_s being the density of the superfluid. Measurements by Clow and Reppy (1966) and Tyson and Douglass (1966) show $\rho_s(T) \sim (T_\lambda - T)^\zeta$ $T < T_\lambda$ with $\zeta \approx 0.666 \pm 0.1$, that is $\beta \sim \frac{1}{3}$. The specific heat at constant pressure shows a logarithmic singularity (Figure 2.3). Experimental results by Fairbank and Kellers (1966), and more accurate results by Ahlers (1969), give $\alpha = 0.000 \pm 0.003$ according to one interpretation of the data and $\alpha = -0.005 \pm 0.005$ according to another.

$$C_P \sim -A \ln \left(1 - \frac{T}{T_\lambda}\right) + B \quad T > T_\lambda \quad . \quad (2.1.20)$$

Henkel, Smith and Reppy (1969) obtained $v_h = 0.67 \pm 0.04$ by measuring the superfluid healing length of thin films which by the assumption of one relevant coherence length near T_λ is the same as v .

Once the Hamiltonian is given the problem is only to evaluate the partition function

$$Z(T, N, V) = \text{Tr } e^{-\beta \mathcal{H}_N} \quad (2.1.21)$$

and then

$$-\frac{F(T,V)}{kT} = \lim_{V \rightarrow \infty} \frac{1}{N} \ln Z \quad (2.1.22)$$

and all thermodynamical quantities follow by the usual relations. The problem of going to the thermodynamic limit in order to have any critical behaviour is at least partially solved by the rigorous proofs of Ruelle (1963) and others who showed that thermodynamical properties are the same whether obtained by canonical or grand canonical ensemble and do not depend on the shape of the sequence of volumes one uses. Still the order in which limits are taken can make a difference and caution is required.

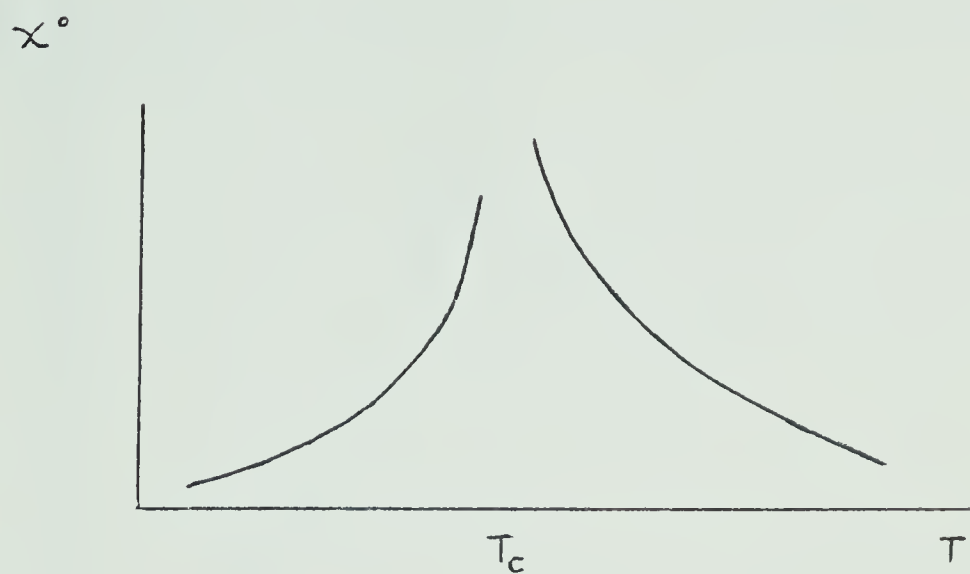


Figure 2.1
Susceptibility in a Ferromagnet

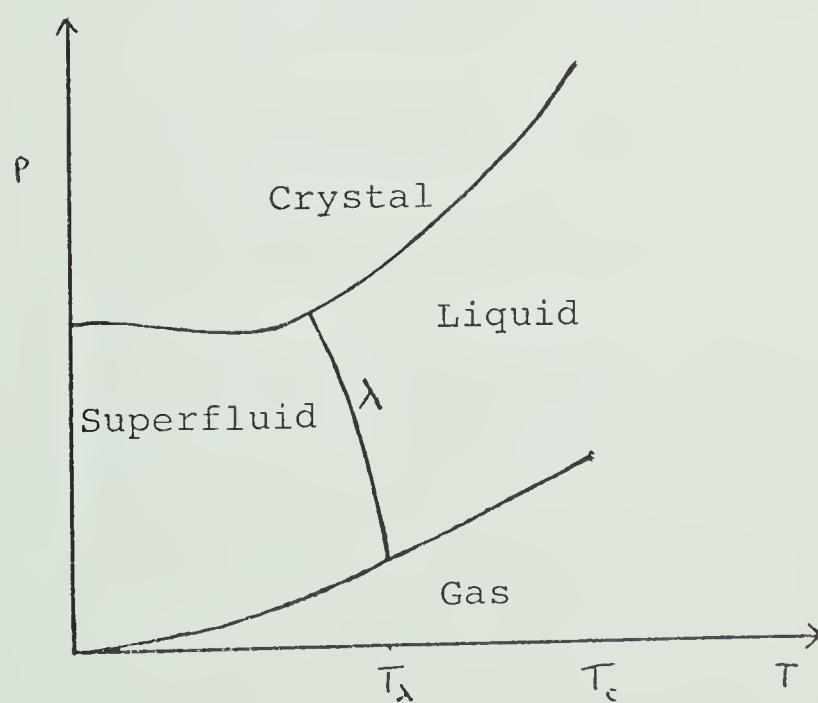


Figure 2.2
Schematic Phase diagram of ^4He

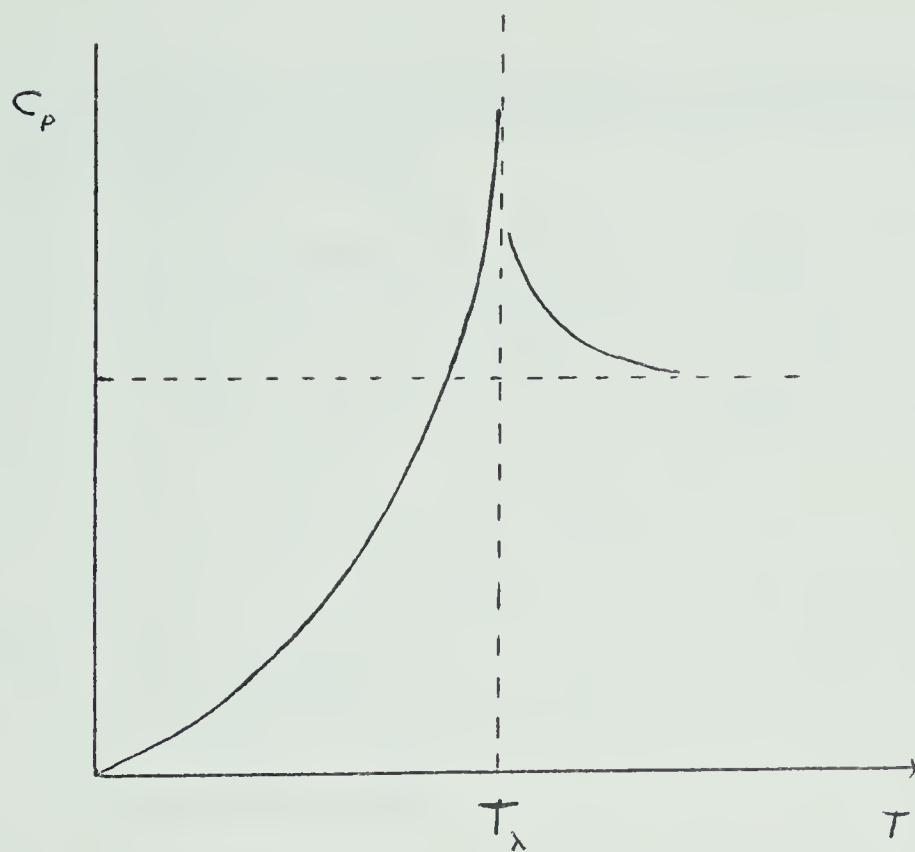


Figure 2.3

Specific heat of liquid helium

2.2 The XY model and its relation to magnetic systems and quantum fluids

The first attempt to build a model for ferromagnetism was the Weiss molecular field theory (1907). The model assumes that each particle in the system interacts equally with all other particles in the system. In spite of this rather unrealistic assumption the results were in some agreement with experiment at least in predicting a critical point, a diverging susceptibility (with $\delta=1$) and a discontinuous specific heat.

The first realistic model to be treated was the Ising model with first neighbour interactions. The spins are fixed on a lattice and the Hamiltonian describing the interaction is

$$\mathcal{H} = -J \sum_{\langle ij \rangle} \sigma_i^z \sigma_j^z - m H^z \sum_i \sigma_i^z \quad (2.2.1)$$

where $\langle ij \rangle$ denotes sum over pairs of nearest neighbours σ_i^z is the Pauli spin matrix and m is the magnetic moment of the atoms.

It was solved for the spin $\frac{1}{2}$ case for a one dimensional chain by Ising (1920), and by Onsager (1944) for a square lattice. Series expansions gave good estimates for the critical indices in 3 dimensional lattices. Magnetic systems were found to agree with the predictions of this model, namely the cobalt tutton salts. The β brass

binary alloy seems to be a good example for this model (Als-Nielsen 1969).

The Heisenberg model is based on the exchange theory of ferromagnetism. The Heisenberg Hamiltonian is

$$\mathcal{H} = -2J \sum_{\langle ij \rangle} \vec{S}_i \cdot \vec{S}_j - m H^Z \sum_i S_i^Z \quad (2.2.2)$$

Materials like EuO and EuS satisfy both assumptions of the model, localization of spins and isotropy of interaction, and series expansion methods gave estimates for the critical indices (Rushbrooke, Wood (1958), Domb, Wood (1964), Baker, Gilbert, Eve and Rushbrooke (1967)).

The natural extension of the Heisenberg model is a model in which the assumption of isotropy is dropped. Mathematically this causes great complications, usually the most general model treated is

$$\mathcal{H} = -2 \sum_{\langle ij \rangle} \{J_{\perp} (S_i^X S_j^X + S_i^Y S_j^Y) + J_{\parallel} S_i^Z S_j^Z\} - m \vec{H} \cdot \sum_i \vec{S}_i \quad (2.2.3)$$

which is called the anisotropic Heisenberg model even though the anisotropy is not complete. The Ising model is the extreme case of $J_{\perp} = 0$. The XY model is the other extreme $J_{\parallel} = 0$. The XY hamiltonian is

$$\mathcal{H}_{xy} = -2 \sum_{\langle ij \rangle} J (S_i^X S_j^X + S_i^Y S_j^Y) \quad (2.2.4)$$

first introduced by Lieb, Schultz and Mattis (1961).

The work done on the anisotropic Heisenberg model was to calculate χ^{zz} which for the limiting case of the XY interaction is not expected to diverge, whereas this thesis deals with a field applied in a direction in the plane of interaction.

A strong inducement to study this model was the work of Jasnow and Wortis on the classical anisotropic Heisenberg hamiltonian. The classical limit is obtained by taking the limit $\vec{S} \rightarrow \infty$ but $\frac{S}{|\vec{S}|} \rightarrow 1$. This gives an interaction of commuting spins which behave like classical vectors. Jasnow and Wortis studied this Hamiltonian for various anisotropies on different lattices by series expansions. Their results suggest, as mentioned in 2.1, that only the symmetry of the ground state and the dimensionality of the lattice affect the critical exponents. In particular as long as the interaction has the symmetry of the XY hamiltonian, that is as long as $|J_{\perp}| > |J_{\parallel}|$, the critical behaviour of the system is that of pure XY interaction. The indices change only when there is a change in the symmetry, and then they change discontinuously. The same result is indicated by Barouch and McCoy (1970) for the one dimensional anisotropic XY model (see 2.3). While the pure XY model is mathematically convenient, it has no possible hope of describing a physical system. However

an anisotropic Heisenberg model with the XY symmetry in which $J_{\perp} \gg J_{\parallel}$ is quite possible.

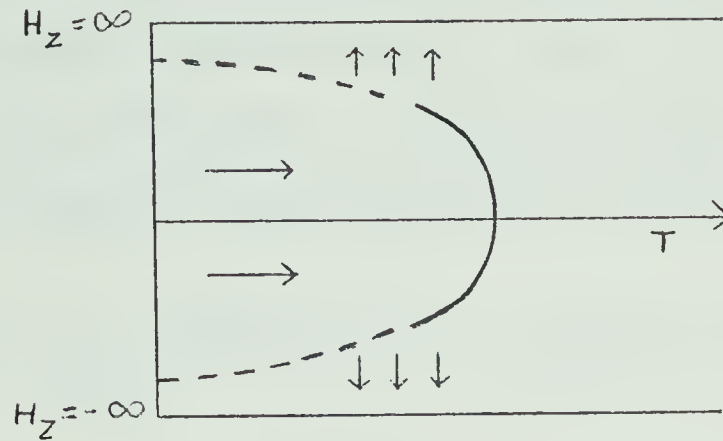
It was suggested by Huiskamp (1967) that materials with strongly anisotropic crystalline fields like $Gd_2(SO_4)_3 \cdot 8H_2O$ would be likely to have an XY like behaviour. It was shown by Betts, Elliott and Lee (1970) that an insulating crystal composed of high half odd spins with a strong axial crystalline field behaves much like a spin $\frac{1}{2}$ XY system. The interaction being

$$\mathcal{H} = + D \sum_i (S_i^z)^2 - J \sum_{\langle ij \rangle} \vec{S}_i \cdot \vec{S}_j \quad D \gg J > 0 \quad (2.2.5)$$

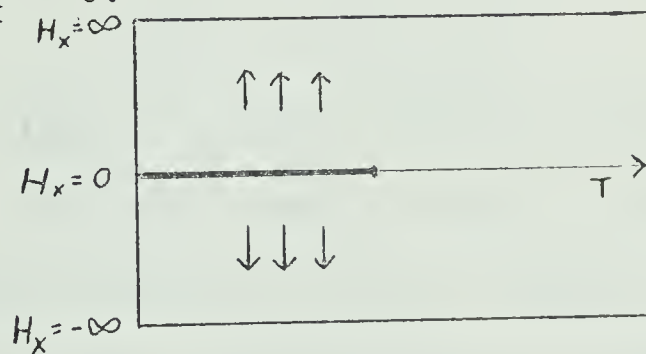
the strong crystalline field will cause the $S^z = \pm \frac{1}{2}$ levels to be highly populated in the critical region. Looking now only at the populated subspace $S^z = \pm \frac{1}{2}$ the ratio of the $S_i^x S_j^x + S_i^y S_j^y$ part of the interaction compared to the $S_i^z S_j^z$ part will be about S^2 to 1.

In such a material we have a dominant transverse coupling, the spontaneous magnetization vector \vec{M} in zero field will lie in the xy plane. Introduction of a field in the z direction will tend to rotate the magnetization towards the z axis for large enough fields and we end with a paramagnetic phase in the xy plane. The expected

phase diagram (Fisher 1968) is



where the broken line indicates the transition may not be second order any more, and we have a λ -like transition line. Introducing a field in the x direction will cause the magnetization to be in the x direction even in the limit of $H_x \rightarrow 0$.



Our calculation is a series expansion for $T > T_c$ $H_x \rightarrow 0$. We approach T_c from above on the $H_x = 0$ line. Using the equalities predicted by scaling, we can obtain the behaviour on other lines, for example the approach to $H_x = 0$ on a $T = T_c$ line.

Matsuda and Matsubara (1956) introduced the anisotropic Heisenberg model as a possible model for a quantum lattice fluid. In the following pages we bring Matsuda and Matsubara's derivation. The model for a classical fluid as

proposed by Yang and Lee (1952) assumes that each atom occupies a cell i with a probability t_i which is either 1 or 0. Multiple occupancy is thus excluded and the grand partition function for the system obtained is:

$$E \propto \sum_{t=0,1} \exp \left[\beta \mu \sum_i t_i - \sum_{ij} \beta \phi(\underline{r}_i - \underline{r}_j) t_i t_j \right] \quad (2.2.6)$$

μ being the chemical potential and $\phi(\underline{r}_i - \underline{r}_j)$ the potential between atoms in cells i and j . Taking

$$\phi(\underline{r}_i - \underline{r}_j) = \begin{cases} 4J & \text{for } i \text{ and } j \text{ nearest neighbours} \\ 0 & \text{otherwise} \end{cases} \quad (2.2.7)$$

the model reduces to the classical Ising model.

For a quantum fluid we work in the framework of second quantization and define creation and annihilation operators a_i^\dagger and a_i respectively on each lattice site i . Since we are dealing with a Bose gas

$$(a_i^\dagger, a_j^\dagger) = (a_i, a_j) = (a_i, a_j^\dagger) = 0 \quad (2.2.8)$$

for $i \neq j$. In order to exclude multiple occupation for $i=j$

$$\{a_i, a_i^\dagger\} = a_i a_i^\dagger + a_i^\dagger a_i = 1 \quad (2.2.9)$$

$$\{a_i^\dagger, a_j^\dagger\} = \{a_i, a_i\} = 0 \quad (2.2.10)$$

and this ensures that the number operators

$$N_i = a_i^\dagger a_i$$

will have eigenvalues 1, 0 only. These commutation relations are satisfied by

$$\begin{aligned} a_i^\dagger &= \sigma_i^X + i \sigma_i^Y \\ a_i &= \sigma_i^X - i \sigma_i^Y . \end{aligned} \quad (2.2.11)$$

The following potential energy ϕ is considered:

$$\left\{ \begin{array}{ll} v = \infty & \text{for two atoms on same lattice point} \\ & \text{(forbidden by (2.2.10))} \\ v = -v_0 & \text{for nearest neighbours} \\ v = 0 & \text{otherwise .} \end{array} \right.$$

Therefore

$$\phi = -v_0 \sum_{ij} a_i^\dagger a_i a_j^\dagger a_j . \quad (2.2.12)$$

As to kinetic energy it is assumed that atoms can move only to vacant nearest neighbour sites. Such a transition is generated by $a_i^\dagger a_j$ for ij nearest neighbours. We look at the form the kinetic energy takes in the continuum

$$\text{K.E.} = \frac{\hbar^2}{2m} \int \frac{\partial \psi^\dagger}{\partial x} \frac{\partial \psi}{\partial x} dx . \quad (2.2.13)$$

For free particles

$$\psi(\underline{r}) = \frac{1}{\sqrt{V}} \sum_{\underline{K}} a_{\underline{K}} e^{-i\underline{K}\underline{r}} \quad (2.2.14)$$

therefore

$$K.E = \frac{\hbar^2 |\underline{K}|^2}{2m} \sum_{\underline{K}} a_{\underline{K}}^\dagger a_{\underline{K}} \quad (2.2.15)$$

if we define for the lattice fluid the operators

$$\begin{aligned} a_{\underline{K}} &= \frac{1}{\sqrt{N}} \sum_{\underline{j}} e^{i\underline{K}\underline{r}_{\underline{j}}} a_{\underline{j}} \\ a_{\underline{K}}^\dagger &= \frac{1}{\sqrt{N}} \sum_{\underline{j}} e^{-i\underline{K}\underline{r}_{\underline{j}}} a_{\underline{j}}^\dagger. \end{aligned} \quad (2.2.16)$$

One possible kinetic energy operator with the correct continuum limit would be

$$K.E = \frac{N\hbar^2 d}{2ma^2 q} \sum_{\langle ij \rangle} (a_i^\dagger - a_j^\dagger)(a_i - a_j) \quad (2.2.17)$$

where we replaced $\frac{\partial}{\partial x}$ by differences, d is the dimensionality, q the number of nearest neighbours

$$\begin{aligned} K.E &= \frac{\hbar^2}{2ma^2 q} \sum_{\langle ij \rangle} \left(\sum_{\underline{K}} (e^{i\underline{K}\underline{r}_i} - e^{i\underline{K}\underline{r}_j}) a_{\underline{K}}^\dagger \right) \times \\ &\quad \times \left(\sum_{\underline{K}'} (e^{-i\underline{K}'\underline{r}_i} - e^{-i\underline{K}'\underline{r}_j}) a_{\underline{K}'} \right) \\ &= \frac{\hbar^2 d}{2ma^2 q} \sum_{\underline{K}\underline{K}'} a_{\underline{K}}^\dagger a_{\underline{K}'} \sum_{\langle ij \rangle} (e^{i\underline{K}\underline{r}_i} - e^{i\underline{K}\underline{r}_j}) \times \\ &\quad \times (e^{-i\underline{K}'\underline{r}_i} - e^{-i\underline{K}'\underline{r}_j}) \end{aligned} \quad (2.2.18)$$

In the limit $|K| \ll \frac{1}{a}$ for ij nearest neighbours

$$e^{i\mathbf{K} \cdot \mathbf{r}_i} - e^{i\mathbf{K} \cdot \mathbf{r}_j} \sim i\mathbf{K} \cdot \mathbf{a} e^{i\mathbf{K} \cdot \mathbf{r}_i}, \quad (2.2.19)$$

therefore

$$\begin{aligned} K.E &\sim \frac{\hbar^2}{2ma^2} \sum_{KK'} a_K^\dagger a_{K'} \sum_{\langle ij \rangle} (\mathbf{K} \cdot \mathbf{a}) e^{i\mathbf{K} \cdot \mathbf{r}_i} e^{-i\mathbf{K}' \cdot \mathbf{r}_i} \\ &\sim \frac{\hbar^2}{2ma^2} \sum_K a_K^\dagger a_K \sum_{\langle ij \rangle} (\mathbf{K} \cdot \mathbf{a})^2. \end{aligned}$$

For a continuum limit the nearest neighbours are isotropically distributed, so $\sum_{\langle ij \rangle} (\mathbf{K} \cdot \mathbf{a})^2 \propto (K)^2 a^2 \frac{q}{d}$

$$K.E \sim \frac{\hbar^2}{2m} |K|^2 \sum_K a_K^\dagger a_K. \quad (2.2.20)$$

The quantum lattice gas Hamiltonian given by this heuristic arguments is

$$\mathcal{H} = \left(\frac{d\hbar^2}{2ma^2q} \right) \sum_{\langle ij \rangle} (a_i^\dagger - a_j^\dagger)(a_i - a_j) - v_0 \sum_{\langle ij \rangle} a_i^\dagger a_i a_j^\dagger a_j. \quad (2.2.21)$$

For spin $\frac{1}{2}$

$$a_i^\dagger a_i = (\sigma_i^x + i\sigma_i^y)(\sigma_i^x - i\sigma_i^y) = 1/2(1 - \sigma_i^z),$$

therefore

$$\mathcal{H} = - \left(\frac{\hbar^2 d}{q 2 m a^2} \right) \sum_{\langle ij \rangle} (\sigma_i^x \sigma_j^x + \sigma_i^y \sigma_j^y) - v_0 \sum_{\langle ij \rangle} \sigma_i^z \sigma_j^z, \quad (2.2.22)$$

which is our anisotropic Heisenberg model. In absence of interactions this reduces to pure XY.

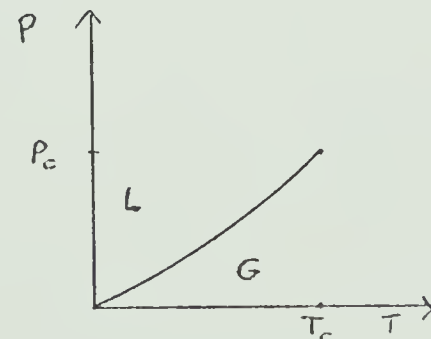
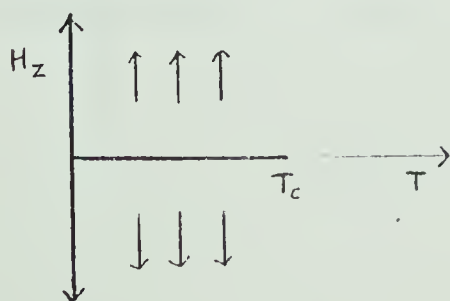
By the assumption that a strong XY interaction compared to the Ising like part yields a behaviour like the pure XY model, we can hope that a pure XY model will describe a quantum fluid with weak interaction be it repulsive or attractive.

This condition that $|J_{\perp}| > |J_{\parallel}|$ is here

$$\frac{d}{v_0 q} \frac{\hbar^2}{2 m a^2} \gg 1, \quad (2.2.23)$$

where a is the lattice spacing, d is the dimensionality, q is the number of nearest neighbours and $-v_0$ the potential energy between two neighbouring atoms.

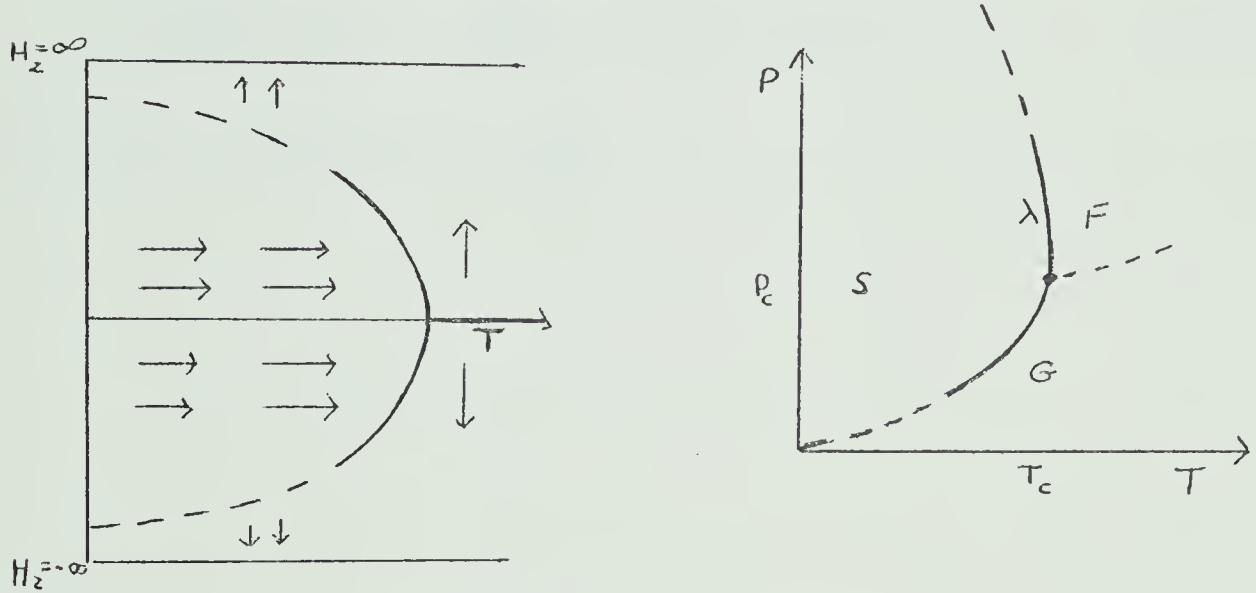
The phase diagram for the superfluid that corresponds to this hamiltonian (2.2.22) can be formed in the same way as the classical gas diagram was formed in analogy to the uniaxial Ising magnet



Our hamiltonian has a positive, ferromagnetic, J as it came from the kinetic energy and is

$$J = \frac{d\hbar^2}{q2ma^2} \quad .$$

The phase diagrams are



The pressure in the fluid still corresponds to a field in the z direction as it is the quantity conjugate to the density, which is $\sum a_i^\dagger a_i \sim \sum \sigma^z$. Applying a field in the z direction would not destroy the order until the field become large enough to rotate the magnetization to the z axis. The magnetic fields in the x or y direction correspond to off diagonal fields, v_i introduced by Bogolyubov (1960), Hohenberg and Martin (1965) which have no physical realization.

$$\mathcal{H}' = - \frac{1}{\sqrt{v_0}} \sum_i (v_i^* a_i + v_i a_i^\dagger) \quad , \quad (2.2.24)$$

then

$$\frac{\partial}{\partial v^*} \frac{kT}{Nv_O} \log Z = \frac{1}{\sqrt{v_O}} \langle a_O \rangle = \langle \psi \rangle . \quad (2.2.25)$$

The superfluid order parameter can be defined by

$$|\psi(T, \mu)| = \lim_{v, v^* \rightarrow 0} |\psi(T, \mu, v^*, v)| . \quad (2.2.26)$$

This is analogous to $\lim_{H_x \rightarrow 0} \langle M^X \rangle$.

2.3 Rigorous one dimensional results

Lieb, Schultz and Mattis (1961) were able to diagonalize the Hamiltonian of the anisotropic XY model in one dimension

$$\mathcal{H} = \sum_{j=1}^N (1+\gamma) S_j^x S_{j+1}^x + (1-\gamma) S_j^y S_{j+1}^y \quad (2.3.1)$$

by using the transformation to fermion operators given by:

$$a_j^\dagger = \exp \left\{ i\pi \sum_{K=1}^{j-1} c_K^\dagger c_K \right\} c_j^\dagger, \quad a_1^\dagger = c_1^\dagger \quad (2.3.2)$$

$$a_j = \exp \left\{ -i\pi \sum_{K=1}^{j-1} c_K^\dagger c_K \right\} c_j, \quad a_1 = c_1. \quad (2.3.3)$$

They calculated the correlation functions $\langle S_0^z S_1^z \rangle$, $\langle S_0^x S_1^x \rangle$ and proved that for the isotropic case, $\gamma=0$, the long range order vanishes even at $T=0$

$$\lim_{n \rightarrow \infty} \langle S_0^x S_{2n}^x \rangle \rightarrow 0, \quad ,$$

while it does not for $\gamma \neq 0$.

Katsura (1962) independently solved the same model. He obtained also the perpendicular susceptibility χ^{zz} in finite fields and the ground state. Suzuki (1966) obtained the staggered initial susceptibility for the isotropic XY model

$$\chi_O = \frac{Nm^2}{kT} \cdot \frac{1}{\pi} \int_0^\pi \frac{d\omega}{\cosh^2(K \cos \omega)} \quad (2.3.4)$$

$$\chi_O^S = \frac{Nm^2}{kT} \cdot \frac{1}{\pi} \int_0^\pi \frac{\tanh(K \cos \omega)}{K \cos \omega} d\omega \quad (2.3.5)$$

where

$$K = \frac{J}{kT} \quad (2.3.6)$$

The ground state functions were obtained by Katsura

$$\psi_O = \left(\prod_{K > N/4}^{N/2} A_{2K-1}^\dagger A_{-2K+1}^\dagger \right) |0\rangle \quad \text{for } \frac{N}{2} \text{ even} \quad (2.3.7)$$

$$\psi_O = \left(\prod_{K > N/4}^{N/2} A_{2K}^\dagger A_{-2K}^\dagger \right) A_N^\dagger |0\rangle \quad \text{for } \frac{N}{2} \text{ odd} \quad (2.3.8)$$

where

$$A_K = \frac{1}{\sqrt{N}} \sum_{\ell=1}^N a_\ell \exp \left[-i\pi \left(\frac{K\ell}{N} - \frac{1}{n} \right) \right] \quad (2.3.9)$$

$$A_K^\dagger = \frac{1}{\sqrt{N}} \sum_{\ell=1}^N a_\ell^\dagger \exp \left[i\pi \left(\frac{K\ell}{N} - \frac{1}{n} \right) \right] \quad (2.3.10)$$

The ground state energy is

$$\frac{E_O}{N|J|} = -\frac{2}{\pi} \quad (2.3.11)$$

These results show that in one dimension there is no long range order. The model behaves like an antiferromagnet with respect to a perpendicular field and the ground state is rather like that of the antiferromagnetic isotropic Heisenberg model.

A detailed study of the correlation functions in the anisotropic XY model was done by Barouch and McCoy (1969, 1970). The hamiltonian treated is:

$$\mathcal{H} = -J \sum [(1+\gamma) S_i^x S_{i+1}^x + (1-\gamma) S_i^y S_{i+1}^y] - \hbar \mu \sum S_i^z. \quad (2.3.12)$$

They studied the asymptotic behaviour of the correlations in finite fields H_z : $\langle S_O^x S_R^x \rangle$, $\langle S_O^y S_R^y \rangle$, and $\langle S_O^z S_R^z \rangle$ as $R \rightarrow \infty$ for high and low temperatures. All correlations decay exponentially with the exponent depending on the field.

For $T=0$ the behaviour is different for the symmetric cases $\gamma=0$, or $\hbar\mu = J$, or $(\hbar\mu/J)^2 + \gamma^2 = 1$. When none of those conditions obtain, all correlations approach their limit exponentially rapidly; when one condition holds they approach the limiting values as some power of R . The effect of those symmetries can be seen at finite T too as the rate of the vanishing depends on the symmetry.

The parallel susceptibility has not been calculated as the solution of the XY chain with a field in the z direction used all simplifying features to the maximum; the nearest neighbour character of the interaction, the one dimensionality and the simple quadratic form of the Hamiltonian. For the parallel susceptibility addition of the appropriate Zeeman term gives the following Hamiltonian:

$$\mathcal{H} = -J \sum S_i^x S_{i+1}^x + S_i^y S_{i+1}^y - m H \sum S_i^x \quad . \quad (2.3.13)$$

Taking $x \rightarrow z$, $z \rightarrow y$, $y \rightarrow x$, we obtain

$$\mathcal{H} = -J \sum (a_i^\dagger a_{i+1} + a_i a_{i+1}^\dagger) - J' \sum a_i^\dagger a_i a_{i+1}^\dagger a_{i+1} + J'' \sum a_i^\dagger a_i \quad (2.3.14)$$

and the simple quadratic nature of \mathcal{H} is destroyed.

High temperature series expansions have been done by Bonner and Fisher (1964). They calculated the energy, specific heat, susceptibility and pair correlations for the hamiltonian

$$\mathcal{H} = -2J \sum [S_i^z S_{i+1}^z + \gamma (S_i^x S_{i+1}^x + S_i^y S_{i+1}^y)] - m \vec{H} \sum \vec{S}_i \quad (2.3.15)$$

for $\gamma=0$ to 1, therefore not reaching the region of XY symmetry.

2.4 Rigorous thermodynamic inequalities and scaling laws

Thermodynamic arguments were used by Rushbrooke (1963, 1965) and Griffiths (1965) to prove rigorous inequalities for the critical exponents. Most of those proofs make use of convexity relations (Ruelle 1963, Fisher 1964 and Griffiths 1964) of the free energy

$$F(\alpha T_1 + \beta T_2, v) \geq \alpha F(T_1, v) + \beta F(T_2, v)$$

$$F(T, \alpha v_1 + \beta v_2) \leq \alpha F(T, v_1) + \beta F(T, v_2) . \quad (2.4.1)$$

The Helmholtz free energy $A(T, v)$ is concave in the temperature and convex in the volume. The same holds with magnetization instead of volume for a hamiltonian of the form

$$\mathcal{H} = \mathcal{H}_0 - H M \quad (2.4.2)$$

or provided the magnetization commutes with the hamiltonian (our case is included in (2.4.2)), we will write the rest in magnetic language.

It follows from the convexity properties that

$$C_M = T \left(\frac{\partial S}{\partial T} \right)_M = -T \left(\frac{\partial^2 A}{\partial T^2} \right)_M \geq 0 \quad (2.4.3)$$

$$\frac{1}{\chi_T} = \left(\frac{\partial H}{\partial M} \right)_T = \left(\frac{\partial^2 A}{\partial M^2} \right)_T \geq 0 \quad (2.4.4)$$

$$C_H = T \left(\frac{\partial S}{\partial T} \right)_H = -T \left(\frac{\partial^2 G}{\partial T^2} \right)_H$$

$$\frac{C_H}{T} = \left(\frac{\partial S}{\partial T} \right)_H = \left(\frac{\partial S}{\partial T} \right)_M + \left(\frac{\partial S}{\partial M} \right)_T \left(\frac{\partial M}{\partial T} \right)_H$$

$$\text{but } \left(\frac{\partial S}{\partial M} \right)_T = - \frac{\partial^2 A}{\partial M \partial T} = - \left(\frac{\partial H}{\partial T} \right)_M = \left(\frac{\partial H}{\partial M} \right)_T \left(\frac{\partial M}{\partial T} \right)_H ,$$

therefore

$$\frac{C_H}{T} = \frac{C_M}{T} + \frac{1}{\chi_T} \left(\frac{\partial M}{\partial T} \right)_H^2 . \quad (2.4.5)$$

Since $C_M \geq 0$

$$C_H \geq T \left(\frac{\partial M}{\partial T} \right)_H^2 / \chi_T \quad T \rightarrow T_C , \quad H = 0 . \quad (2.4.6)$$

Substituting the expected behaviour near the critical point from below at zero field

$$C_H \sim (T_C - T)^{-\alpha} ,$$

$$\chi_T \sim (T_C - T)^{-\gamma} ,$$

$$\left(\frac{\partial M}{\partial T} \right)_H \sim (T_C - T)^{\beta-1} ,$$

we obtain

$$f(T) (T_C - T)^{-\alpha} \geq (T_C - T)^{2(\beta-1)+\gamma} g(T) ,$$

where f and g are non-singular at T_C . Since this should hold for all T close to T_C , it follows that

$$\alpha' + 2\beta + \gamma' \geq 2 \quad . \quad (2.4.7)$$

This is usually called the Rushbrooke inequality.

The implicit assumption in the above derivation was the existence of γ' , that is

$$\lim_{H \rightarrow 0} \chi_T(H, T) = \infty \quad \text{for } T \rightarrow T_C^- .$$

The same assumption and use of convexity of $A(T, M)$ yield (Griffiths 1965)

$$\alpha' + \beta(1 + \delta) > 2 \quad . \quad (2.4.8)$$

Many other inequalities were obtained such as

$$\gamma' \geq \beta(\delta - 1) \quad (2.4.9)$$

$$\gamma(\delta + 1) \geq (2 - \alpha)(\delta - 1) \quad . \quad (2.4.10)$$

The assumptions made to obtain them are plausible for most systems, such as $\alpha' \leq \alpha$ or

$$\left(\frac{\partial M^2}{\partial H^2} \right)_T \leq 0 \quad \text{for } H \geq 0 \quad .$$

Buckingham and Gunton (1969) and Fisher (1969) have proved the following inequalities

$$(2 - \eta)\nu \geq \gamma \quad (2.4.11)$$

$$d \frac{\delta - 1}{\delta + 1} \geq 2 - \eta \quad (2.4.12)$$

and

$$\frac{d\gamma'}{2\beta+\gamma'} \geq 2 - \eta \quad (2.4.13)$$

where d is the dimensionality of the system.

The additional assumptions that went into those two inequalities 12, 13 are that for all temperatures and $H \geq 0$

$$\begin{aligned} \Gamma_1 &\geq 0 \\ \Gamma_2(r) - \Gamma_1^2 &\geq 0 \\ \Gamma_4(r_1, r_2, r_3) - \Gamma_2(r_1) \Gamma_2(r_3 - r_2) &\geq 0 \end{aligned} \quad (2.4.14)$$

where Γ_2 is the 2 spin correlation function and therefore Γ_1 is proportional to the spontaneous magnetization,

$$\Gamma_1(T, H), \Gamma_2(T, H) \text{ and } \Gamma_4(T, H) \quad (2.4.15)$$

are monotonic non-decreasing functions of H ,

$$\Gamma_1(T, H), \Gamma_2(T, H) \text{ and } \Gamma_4(T, H) \quad (2.4.16)$$

are monotonic non-increasing functions of T for any field $H \geq 0$. These assumptions have been proved for the Ising model (Griffiths (1967, 1970)).

The Josephson inequalities (1967) are

$$d\nu' > 2 - \alpha' \quad (2.4.17)$$

$$d\nu > 2 - \alpha \quad .$$

His assumptions are very general such as constant volume or constant pressure constraints.

One way of stating the basic assumption of static scaling theory (Griffiths (1968)) is to require that the singular part of the Helmholtz potential \hat{A} be a homogeneous function of ϵ and $|M|^{1/\beta}$

$$\hat{A}(\lambda\epsilon|\lambda^\beta M) = \lambda^p \hat{A}(\epsilon, M) \quad (2.4.18)$$

p is an arbitrary scaling parameter and $\epsilon = (T - T_c)/T_c$.

This means that, over the M, ϵ plane, the singular part of A in an area around the origin will be the same as for a larger scaled area when itself properly scaled. This was first proposed by Widom (1965) and given physical meaning by Kadanoff (1966) who derived this homogeneity for the Ising model by dividing the system into cells of L lattice sites per side, with L large but smaller than $\frac{\xi}{a}$ (ξ : the coherence length, a is the lattice spacing), obtaining the free energy and claiming that it should not depend on L .

Taking first derivative of (2.4.18) with respect to M and substituting

$$\hat{H} = - \left(\frac{\partial \hat{A}}{\partial M} \right)_T$$

we have

$$\lambda^\beta \hat{H}(\lambda \varepsilon, \lambda^\beta M) = \lambda^p \hat{H}(\varepsilon, M) \quad . \quad (2.4.19)$$

Taking a second derivative with respect to M and using

$$\frac{1}{\hat{\chi}_T} = \left(\frac{\partial H}{\partial M} \right)_T , \text{ we obtain}$$

$$\lambda^\beta \hat{\chi}_T^{-1}(\lambda \varepsilon, \lambda^\beta M) = \lambda^p \hat{\chi}_T^{-1}(\varepsilon, M) \quad . \quad (2.4.20)$$

On the critical isotherm $H \sim M^\delta$, setting that in (2.1.20) equating powers of λ

$$\beta + \beta\delta = p \quad .$$

Setting $M=0$ and $T \rightarrow T_c$ in (2.1.18) yields

$$2\beta + \gamma' = p \quad . \quad (2.4.21)$$

For $M=0$, $T \rightarrow T_c^+$ we obtain $2\beta + \gamma = p$. Eliminating p gives the Widom relation

$$\gamma' = \beta(\delta - 1) \quad (2.4.22)$$

and

$$\gamma = \gamma' \quad . \quad (2.4.23)$$

Taking more derivatives yields more such relations.

All the rigorous inequalities we quoted in the first part of this section become equalities under the scaling hypothesis (there are others we did not mention that do not). Since A is a function of two variables we can see

that only two independent critical exponents are allowed by the scaling hypothesis.

Taking the derivatives of the free energy and using the same homogeneity argument one has a relation for the gap exponents defined in (2.1.7)

$$F^{(3)}(\epsilon) = \lambda^{2p-3\beta} F^{(3)}(\lambda\epsilon)$$

but

$$F^{(3)}(\epsilon) = (-\epsilon)^{-\Delta'_3} F^2(\epsilon) \sim (-\epsilon)^{-\Delta'_3-\gamma'}$$

so

$$\Delta'_3 = 2p - 3\beta - \gamma' = \beta + \gamma'$$

but

$$\Delta'_2 = 2 - \alpha' - \beta = \beta + \gamma'$$

and

$$\Delta'_2 = \beta + \gamma' \quad .$$

Similarly it can be shown that all gap parameters for $T < T_c$ and all even ones for $T > T_c$ are equal.

These results of scaling hold exactly in the 2-dimensional Ising model, and are approximately true for 3-dimensional Ising and Heisenberg models. Some discrepancies can be explained by allowing larger errors in the numerical calculations.

The Kadanoff construction applied to correlation functions yields a scaling behaviour for them which already for the 3-dimensional Ising model is outside the range predicted by numerical calculations.

Since these arguments are extendable to dynamical scaling we will deal with the correlation scaling in the dynamical section of this thesis, and mention here only that the equalities obtained for the static indices are (2.4.11) and (2.4.17) with the equality holding, that is:

$$\begin{aligned}
 dv' &= 2 - \alpha' & dv &= 2 - \alpha \\
 (2-\eta)v' &= \gamma' & (2-\eta)v &= \gamma & (2.4.24) \\
 d \frac{\delta-1}{\delta+1} &= \frac{d\gamma'}{2\beta+\gamma'} = 2 - \eta .
 \end{aligned}$$

When we assume scaling all those equalities will be assumed.

CHAPTER III

HIGH TEMPERATURE SERIES EXPANSIONS

3.1 General method of graphical expansions

The free energy can be expressed as a formal series in powers of T^{-1} . The method has been expounded by Domb (1960,1965) and by Fisher (1965) in their review articles. We have $Z = \text{Tr} \exp((-H/kT)$ and

$$-\frac{F}{kT} = \lim_{N \rightarrow \infty} \frac{1}{N} (\log Z), \text{ therefore}$$

$$-\frac{F}{kT} = \lim_{N \rightarrow \infty} \frac{1}{N} \log \sum_{n=0}^{\infty} \frac{K^n}{n!} \text{Tr } P^n \quad (3.1.1)$$

where $K = J/kT$ and P is defined by $H = -JP$. As shown by Rushbrooke and Wood (1958) the logarithm has the effect of taking only the terms linear in N . The free energy being an extensive variable this was to be expected.

$$\log Z = 2 \log N + \log \left(1 + \sum_{n=1}^{\infty} \mu_n \frac{K^n}{n!} \right) \quad (3.1.2)$$

where $\mu_n = \frac{\text{Tr } P^n}{\text{Tr } I}$. For the spin $\frac{1}{2}$ case $\text{Tr } I = 2^N$, I being the direct product of N spinor unit matrices.

The coefficients μ_n are polynomials in N and we assume that the logarithm is also expandable in a power series in K

$$\log \left(1 + \sum_{n=1}^{\infty} \mu_n \frac{K^n}{n!} \right) = \sum_{n=1}^{\infty} \lambda_n \frac{K^n}{n!} \quad (3.1.3)$$

which is a definition of the coefficients λ_n . The result can be checked using

$$\lambda_n = \sum_{m=1}^n \frac{(-1)^{m+1}}{m} \mu_n^{(m)}$$

but has been proved generally. Hence

$$-\frac{F}{kT} = \sum_{n=0}^{\infty} \frac{K^n}{n!} \text{Tr}^* P^n \quad (3.1.4)$$

where $\text{Tr}^* A = \frac{(\text{part linear in } N \text{ of } \text{Tr } A)}{\text{Tr } I}$.

Our hamiltonian, like all hamiltonians used in critical phenomena, is a sum of pairwise interactions; therefore each term in $\text{Tr}^* P^n$ can be represented diagrammatically as a linear graph of n lines and $v \leq 2n$ vertices.

$$\text{Tr}^* P^n = \sum_i w(g_i^n) (g_i^n, L) \quad (3.1.5)$$

g_i^n are all possible linear graphs of n lines (i is a dummy index ordering them). (g_i^n, L) is the lattice constant of the graph g_i^n determined by the number of ways the graph g_i^n can be embedded in a lattice L . $w(g_i^n)$ is a weight assigned to the graph, independently of the lattice, depending on the interaction and the graph.

3.2 The graphical expansion method used for the XY model

The properties of the operators a_i and a_i^\dagger simplify the calculations. The method used by Betts, Elliott and Lee (1970) is as follows: The general term wanted for the free energy is $\text{Tr}^* P^n$. Each such term will be represented by graphs with n arrows as each power of P contributes a factor $a_i^\dagger a_j$ which can be seen as an arrow from i to j . Therefore a priori all linear graphs of n arrows, and v vertices $v \leq 2n$, can be expected to appear in $\text{Tr}^* P^n$.

A non zero contribution can only be obtained when:
 (a) an equal number of arrows enter and leave each site;
 (b) arrow heads and tails alternate at each site.

This follows because

$$a_i a_i = a_i^\dagger a_i^\dagger = 0 \quad (3.2.1)$$

and

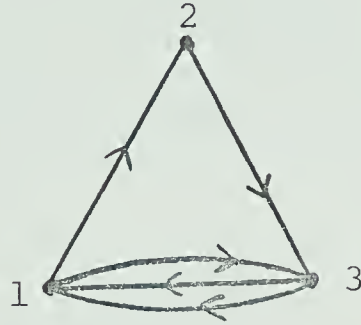
$$\text{Tr } a_i = \text{Tr } a_i^\dagger = 0 . \quad (3.2.2)$$

Were there an odd number of operators at a site we could reduce them, two at a time using $\{a_i, a_i^\dagger\} = 1$, to a single traceless operator.

The partition function is the sum of all ordered directed graphs satisfying (a). Several of those will correspond to one directed shadow graph. Several shadow graphs will correspond to one bare graph in which bonds replace the directed arrows.

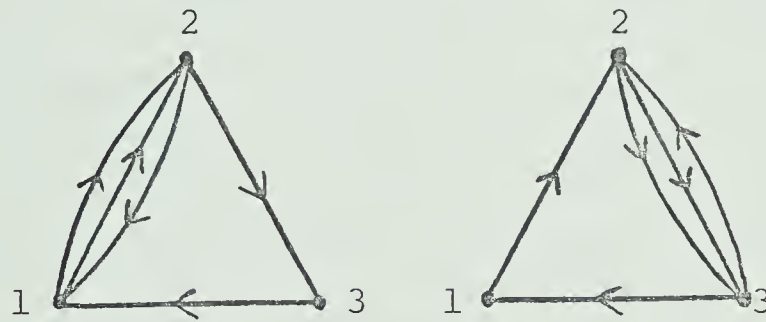
Example

The term $\text{Tr}^* a_1^\dagger a_2^\dagger a_3^\dagger a_1^\dagger a_3^\dagger a_2^\dagger a_1^\dagger$ corresponds to the shadow graph



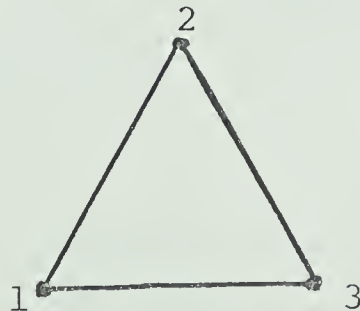
Out of $5!=120$ possible orderings of the 5 arrows only 20 will give a non zero contribution and they all correspond to the same shadow graph.

This shadow graph is equivalent to 5 other shadow graphs, the following two



and those obtained from these three by reversing all arrows.

All six are represented by the same bare graph



The trace of each such graph is 2^{N-s} , where s is the number of vertices of the bare graph, since for spin $\frac{1}{2}$ the trace of the unit matrix is 2 while $\text{Tr} a_i^\dagger a_i =$

$\text{Tr } a_i a_i^\dagger = 1$. This is a great simplification compared to the Heisenberg model where traces of different spin operators have to be calculated.

The horizontal weight of each shadow graph g' is defined by

$$h(g') = \varepsilon(g') S(g)/S(g') \quad (3.2.3)$$

where $S(g)$ is the symmetry number of the bare graph g , $S(g')$ is the symmetry number of the shadow directed graph g' and $\varepsilon(g')$ equals one if g' is equivalent to the graph obtained from it by reversing all arrows. Going back to the above example, the symmetry number of the triangle is 6, and the symmetry number of our graph is 1. $\varepsilon(g') = 1$ as in this case interchange of two labels is the same as reversal of arrows. Equivalence of shadow graphs is defined as follows: two graphs are equivalent if their arrows and vertices are in one to one correspondence. Such a correspondence is an embedding. The number of embeddings of one graph in itself is its symmetry number.

The vertical weight $v(g')$ is defined as the number of allowed ordered directed graphs corresponding to the shadow graph g' . In our example that number was 20. Allowed here meant giving a non zero contribution, therefore satisfying criterion (b).

All these weights are calculated by computer; for that purpose graphs are represented, as usual, by matrices or vectors. For the directed shadow graph the matrix element $M(i,j)$ is the number of arrows from i to j . The bare graph matrix has a unit element in a location i,j if the vertices i and j are connected, and zero otherwise. For our example

$$M(g') = \begin{pmatrix} 0 & 1 & 1 \\ 0 & 0 & 1 \\ 2 & 0 & 0 \end{pmatrix}, \quad M(g) = \begin{pmatrix} 0 & 1 & 1 \\ 1 & 0 & 1 \\ 1 & 1 & 0 \end{pmatrix}.$$

In order to exclude equivalent graphs, to find among a group of lattice constants the one pertaining to the bare graph and to calculate the horizontal weight the incidence matrices defined above were used.

For the vertical weight another kind of matrix representation was used. The matrix is rectangular. Rows denote order of arrows, and the columns specify vertices. The arrow in the k 'th position being i - j the matrix elements in the k 'th row will be $N(k,i)=1$, $N(k,j)=-1$. The vertical weight of the graph corresponds therefore to the number of permutations of rows of the matrix N that alternate 1 and -1 in each column. Full use of efficient methods of counting permutations was made in the computer program written by Dr. C.J. Elliott for the free energy calculation. This calculation being rather

time consuming for higher order graphs, several theorems were proved (Betts, Elliott and Lee (1970)) which reduced noticeably the computer time required and were based on full use of the fact that there are no restrictions at vertices of order 2. We will mention those used and their generalizations for our calculations in the next section.

To summarize, the final contribution to the term $\text{Tr}^* P^n$ is

$$\sum_{g'_{n,i}} (g_{n,i}, L) v(g'_{n,i}) h(g'_{n,i}) 2^{-s} \quad (3.2.4)$$

where the sum is over all inequivalent directed shadow graphs of n arrows, $g'_{n,i}$. Note that the lattice constant is that of the bare graph $g_{n,i}$ the horizontal weight giving the occurrence of the directed graph $g'_{n,i}$ for each bare graph, the vertical weight being the number of allowed orders of the arrows and 2^{-s} being the actual trace.

3.3 The series expansion method for the initial parallel susceptibility

A. The XY model is invariant under rotation in the xy plane, therefore we can choose any direction in that plane to calculate the initial parallel susceptibility. We add a Zeeman term and the resulting hamiltonian is:

$$\mathcal{H} = -J \sum_{\langle i,j \rangle} (a_i^+ a_j + a_i a_j^+) - H_x M_x \quad . \quad (3.3.1)$$

The initial parallel susceptibility is defined by

$$\chi_{||}^0 \equiv \chi_o^{xx} = \frac{1}{N\beta} \left. \frac{\partial^2}{\partial H_x^2} \log Z \right|_{H_x=0} \quad . \quad (3.3.2)$$

For models where $[M, H] = 0$ this reduces to $\frac{1}{N} \langle M_x^2 \rangle$, which is the square of the fluctuation in the long range order. The relation between the two is discussed in Subsection 3.3B. Using the notation of the last section, $\mathcal{H} = -J P$, we have:

$$\chi_{||}^0 = \frac{1}{N\beta} \left. \frac{\partial^2}{\partial H_x^2} \log \text{Tr} e^{\beta(JP + H_x M_x)} \right|_{H_x=0} \quad . \quad (3.3.3)$$

As before we use the relation

$$\frac{1}{N} \log \text{Tr} e^{\beta A} = \sum \frac{\beta^n}{n!} \text{Tr}^* A^n \quad (3.3.4)$$

and obtain:

$$\chi_{\parallel}^0 = \frac{1}{\beta} \frac{\partial^2}{\partial H_X^2} \sum \frac{\beta^n}{n!} \text{Tr}^* (JP + H_X M_X)^n \Big|_{H_X=0} . \quad (3.3.5)$$

We take the second derivative and set $H_X=0$. This means keeping only the term proportional to H_X^2 . The coefficient of H_X^2 in $\text{Tr}^* (JP + H_X M_X)^n$ can be written in the following form:

$$\frac{1}{2} \text{Tr}^* \sum_{k=0}^{n-2} n(JP)^k M_X (JP)^{n-2-k} M_X . \quad (3.3.6)$$

We obtain (3.3.6) as follows: The term with two factors B in $\text{Tr} (A + B)^n$ is

$$\text{Tr} \sum_{\ell=0}^{n-2} \sum_{s=0}^{n-2} A^\ell B A^{n-2-\ell-s} B A^s .$$

We use the cyclic property of the trace to bring the factor A^s to the front. The last expression reduces to

$$\text{Tr} \sum_{j=0}^{n-2} (j+1) A^j B A^{n-2-j} B \quad (3.3.7)$$

as $j = \ell+s$ can be obtained in $j+1$ ways by adding two integers ℓ and s . Again we invoke the cyclic property of the trace and note that $\text{Tr} A^j B A^{n-2-j} B = \text{Tr} A^{j'} B A^{n-2-j'} B$ for $j' = n-2-j$. We add those two terms and divide by two,

$$\begin{aligned}
& \frac{1}{2} [(j+1) + (n-2-j+1)] \operatorname{Tr} A^j B A^{n-2-j} B \\
& = \frac{1}{2} n \operatorname{Tr} A^j B A^{n-2-j} B .
\end{aligned} \tag{3.3.8}$$

For even n the sum (3.3.7) has an odd number of terms. Still the expression (3.3.8) holds because the middle term occurs for $j = \frac{n-2}{2}$ and is $\frac{n}{2} \operatorname{Tr} A^j B A^{n-2-j} B$.

Now that we have justified the use of (3.3.6) we can substitute it in (3.3.5) and obtain

$$\chi_{II}^O = \frac{1}{J} \sum_n \frac{K^{n-1}}{(n-1)!} \sum_{k=0}^{n-2} \operatorname{Tr}^* P^k M_x P^{n-2-k} M_x . \tag{3.3.9}$$

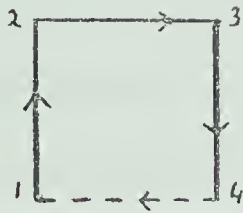
We have $M_x = \frac{m}{2} \sum (a_i^\dagger + a_i)$, hence

$$\chi_{II}^O = \frac{1}{J} \sum_n \frac{K^n}{n!} \sum_{k=0}^{n-1} \sum_{i,j} \operatorname{Tr}^* (P^k a_i^\dagger P^{n-1-k} a_j + P^k a_i P^{n-1-k} a_j^\dagger) . \tag{3.3.10}$$

Each term in $\operatorname{Tr}^* (P^k a_i^\dagger P^{n-1-k} a_j)$ can be represented by a graph with $n-1$ solid arrows and a dotted arrow. The dotted arrow represents the term $a_i^\dagger a_j$ which is not restricted to nearest neighbours. For $i=j$ we have a graph of $n-1$ solid arrows and a dotted "bubble" which can be on a vertex of the graph or off the graph.

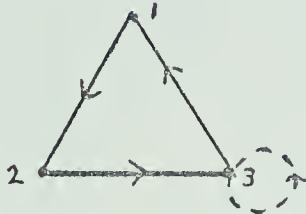
The calculation is done separately for the three cases:

(a) $i \neq j$ Each ij term occurs twice therefore the extra factor 2 in (3.3.12).



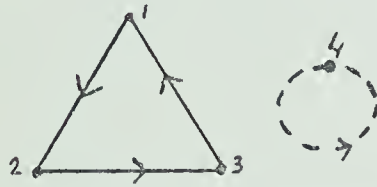
$$a_1^\dagger a_2 a_2^\dagger a_3 a_3^\dagger a_4 a_4^\dagger a_1$$

(b) $i=j$ and i is on the graph



$$a_1^\dagger a_2 a_3 a_2^\dagger a_3^\dagger a_3 a_1 a_3$$

(c) $i=j$ and i is not on the graph



$$a_1^\dagger a_2 a_2^\dagger a_3 a_3^\dagger a_1 a_4^\dagger a_4$$

Case (c) is the simplest. Since i is not on the graph g_{n-1} , a_i^\dagger and a_i commute with the factors P and

$$\text{Tr}^* (P^k a_i^\dagger P^{n-1-k} a_i + P^k a_i P^{n-1-k} a_i^\dagger) = \text{Tr}^* P^{n-1}.$$

But here $\text{Tr}^* P^{n-1}$ does not give the same contribution as g_{n-1} gave in the partition function. The lattice constant is now the separated lattice constant $(g_{n-1}; L)$, which is $-s(g_{n-1}; L)$. s is the number of vertices of g_{n-1} .

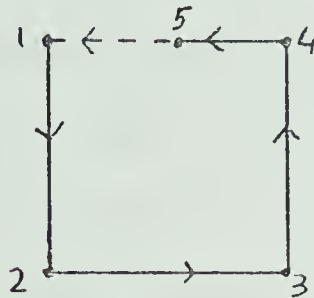
The other factors, horizontal and vertical weights and trace are the same as in the partition function.

The resulting series is:

$$\chi_{||}^{o(c)} = - \frac{K}{J} \sum \frac{K^{n-1}}{(n-1)!} \sum_{g_{n-1}} \frac{h(g_{n-1})(g_{n-1}; L) s V(g_{n-1})}{2^s} \quad (3.3.11)$$

The weights are given in Table A.9 in the Appendix.

For case (a) $i \neq j$, the dotted arrow is placed with the head at the last level and we sum over all allowed levels for its tail. Since vertices of order 2 add no restrictions we introduce, for the vertical weight calculation only, a spurious vertex in the dotted arrow.



$$a_1^\dagger a_2 a_2^\dagger a_3 a_4^\dagger a_5 a_3^\dagger a_4 a_5^\dagger a_1$$

We allow the new arrow formed by the former tail and the spurious vertex to permute freely with the other solid arrows. The arrow formed by the spurious vertex and the tail is kept fixed at the last level.

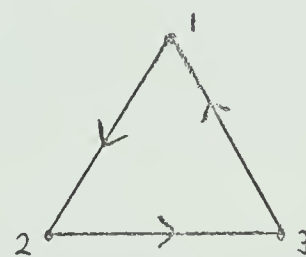
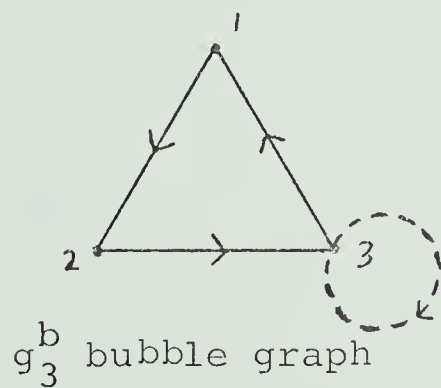
The lattice constant is calculated after deleting the dotted arrow since it is not restricted to nearest neighbours. The horizontal weight is as defined in (3.2.3). In order to keep the dotted arrow different from the others we put a minus sign in front of the matrix element $M(i,j)$. $N(i,j)$ and $N(j,i)$ become zero if no other arrow connects i and j .

The graphs were obtained by Betts, Elliott and Lee (1970) for the fluctuation series. We calculated the new vertical weights only

$$\chi_{||}^{o(a)} = \frac{2}{J} \sum_n \frac{K^n}{n!} \sum_{g_{n-1}''} \frac{h(g_{n-1}'') v(g_{n-1}'') (g_{n-1}, L)}{2^S} \quad (3.3.12)$$

where g_{n-1}'' is the graph with $n-1$ solid arrows and a dotted arrow. g_{n-1} is the bare graph with the dotted arrow removed. The weights are given in Table A.5.

In case (b) we have $i = j$, $i \in g$. The same procedure as for case (a) is followed for the vertical weight calculation. Lattice constants are those of the bare graph without the bubble. The graphs were generated by taking the partition function graph adding a "bubble" at each vertex and eliminating equivalent graphs. We kept the spurious vertex, $s+1$, in the generation process and the horizontal weight calculation. In order to keep the bubble different we introduced a minus sign in $M(i, s+1)$, $M(s+1, i)$ and set $N(i, s+1) = N(s+1, i) = 0$, that is the bubble is deleted in the bare graph. The result is reflected in the following example



Let g_{n-1}^b be a partition function graph g_{n-1} with a bubble on, then:

$$\chi_{\parallel}^o(b) = \frac{1}{J} \sum \frac{K^n}{n!} \sum_{g_{n-1}^b} \frac{h(g_{n-1}^b) v(g_{n-1}^b) (g_{n-1}, L)}{2^S} \quad (3.3.13)$$

The weights can be calculated from Table A.7.

The calculation of fluctuation type vertical weights (one arrow fixed) is facilitated by the lemmas and theorems of Betts, Elliott and Lee (1970).

Lemma 1

Let g be a graph with ℓ arrows and g^* is the graph obtained from g by turning one arrow into a dotted arrow. Let there be ℓ_1 arrows from vertex a to vertex b one of which is the dotted arrow. Then

$$v_{g^*} = \frac{v_g \cdot \ell_1}{\ell} \quad .$$

Proof

If $\ell_1 = 1$ fixing one arrow at a level reduces the number of possible vertical orderings by $\frac{1}{\ell}$. The reason is in the cyclical property of the trace, all cyclic permutations of an allowed permutation are allowed.

If $\ell_1 > 1$ there are ℓ_1 arrows which were indistinguishable in g and by fixing one we made them distinguishable.

Lemma 2

Let g be a graph with ℓ arrows and s vertices and let one arrow be of order 2 in both vertices. Then, one vertex of order 2 can be eliminated. The resulting graph of $\ell-1$ arrows and $s-1$ vertices has a vertical weight v_1 .

$$v = v_1 \cdot \ell \quad .$$

The procedure can be repeated and all arrows of order 2 in both vertices eliminated. The proof is obvious.

Theorem I

Let g be a disconnected shadow graph of ℓ arrows consisting of two subgraphs g_1 and g_2 having no vertices in common. If g_1 and g_2 , consisting of ℓ_1 and ℓ_2 arrows respectively, have vertical weights v_1 and v_2 , then the vertical weight, v , of g is given by

$$v = (\ell! / \ell_1! \ell_2!) v_1 v_2 \quad . \quad (3.3.14)$$

Proof

As there is no interference between the arrangements in the two subgraphs the vertical weight of the graph g is given by the product of two factors. The first is the number of ways of assigning ℓ vertical levels such that ℓ_1 are occupied by the arrows of g_1 and the remaining $\ell_2 = \ell - \ell_1$ are occupied by the arrows of g_2 . The other factor is then the product of the vertical weights of the component subgraphs.

Theorem II

Let a connected shadow graph g of ℓ arrows be composed of two subgraphs g_1 and g_2 of ℓ_1 and ℓ_2 arrows respectively such that g_1 and g_2 have m common vertices of order 2 in g (hence of order 1 in each of g_1 and g_2)

and no other common vertices. Then the vertical weight, v , of g is given by the same expression as in theorem I.

Proof

As there is no restriction on head and tail vertical sequence at vertices of order 2, the vertical weight of g must be the same as that of a separated graph g' consisting of subgraphs g_1 and g_2 . But the vertical weight of g' is given by (3.3.14) and so the theorem is established.

Theorem III

Consider a shadow graph g of order ℓ composed of two subgraphs g_1 and g_2 of $\ell_1 - 1$ and $\ell_2 - 1$ arrows respectively and a third subgraph g_3 consisting of a single arrow. The subgraphs g_1 and g_2 have in common m vertices of order 2 in g and no other common vertices. The subgraph g_3 (the single arrow) connects vertices of arbitrary (even) order orders r_1 and r_2 in g_1 and g_2 respectively. Then the vertical weight of g is given by (3.3.14) again. However, now v_1 and v_2 are the vertical weights of the graphs g'_1 and g'_2 where g'_1 is a subgraph of g consisting of g_1 and g_3 , and g'_2 is a subgraph consisting of g_2 and g_3 . It is important also to note that $\ell_1 + \ell_2 = \ell + 1$.

Proof

Observe first that the vertical weight of g is the same as that of the graph g' obtained from g by separating

all the second order vertices. Next observe that in counting the number of allowed arrow configurations for any shadow graph of order ℓ we may assume that one arrow is at a fixed level and compute the number, n , of allowed configurations of the remaining arrows. Then the vertical weight, $v=\ell n$.

In g' let the arrow forming g_3 be fixed. The remaining $\ell - 1$ arrows belong either to g_1 or g_2 . The $\ell_1 - 1$ arrows may be arranged among the $\ell_1 - 1$ levels assigned to g_1 in n_1 ways independently of the arrangement in n_2 ways of the $\ell_2 - 1$ arrows of g_2 among the $\ell_2 - 1$ levels assigned to g_2 . The $\ell - 1$ levels may be assigned in $(\ell - 1)!/(\ell_1 - 1)!(\ell_2 - 1)$ ways giving

$$n = n_1 n_2 (\ell - 1)!/(\ell_1 - 1)!(\ell_2 - 1) \quad . \quad (3.3.15)$$

Noting that $n_1 = v_1/\ell_1$ and $n_2 = v_2/\ell_2$ the result (3.3.14) follows.

B. We discuss here the relation of the fluctuation to the initial susceptibility.

As seen in Sec. 5.4

$$\chi_T = \int_0^\beta \langle M e^{\lambda \mathcal{H}} M e^{-\lambda \mathcal{H}} \rangle d\lambda \quad (3.3.16)$$

but

$$e^{\lambda \mathcal{H}} M e^{-\lambda \mathcal{H}} = M + \lambda [\mathcal{H}, M] + \frac{\lambda^2}{2!} [\mathcal{H}, [\mathcal{H}, M]] + \dots \quad (3.3.17)$$

therefore

$$\chi_T = \beta Y + \frac{\beta^2}{2!} \langle M [\mathcal{H}, M] \rangle + \dots \quad (3.3.18)$$

Each successive term in (3.3.18), when expanded in a high temperature series, starts its contribution at a higher power of β and its contribution is less divergent. The reason is that in a term $\langle M [\mathcal{H}', [\mathcal{H}, [\dots [\mathcal{H}, M]]] \dots] \rangle$ each commutator has the effect of removing two spins.

Falk and Bruch (1969) proved that for a general thermodynamical system

$$\frac{1 - e^{-\frac{1}{2}\beta\bar{\omega}}}{\frac{1}{2}\beta\bar{\omega}} \leq \frac{\chi_T}{\beta Y} \leq 1 \quad (3.3.19)$$

where

$$\bar{\omega} = \frac{\langle [M, [\mathcal{H}, M]] \rangle}{2Y} \quad (3.3.20)$$

The equality of the indices for Y and χ_T follows from the assumption that the numerator in $\bar{\omega}$ is finite as $T \rightarrow T_c$. Though $\langle [M, [\mathcal{H}, M]] \rangle$ has been shown to be finite for the isotropic Heisenberg model by Mermin and Wagner (1966), their proof cannot be generalized to the parallel susceptibility of the XY model.

Our results indicate that $\langle [M, [\mathcal{H}, M]] \rangle$, which is the first moment of the frequency dependent susceptibility

is not finite. As $T \rightarrow T_c$ the first moment diverges with exponent $\alpha_2 \approx 0.09$. Nevertheless $\omega \xrightarrow{T \rightarrow T_c} 0$ because the fluctuation, Y , diverges much more strongly.

The proof of Falk and Bruch follows from the following expressions for the susceptibility

$$\chi_T = 2 \int_0^{\beta/2} d\lambda \langle e^{\lambda \mathcal{H}_M} e^{-\lambda \mathcal{H}_M} \rangle \quad (3.3.21)$$

and

$$\chi_T = \int_{-\infty}^{\infty} d\omega S_O(\omega) \frac{(1 - e^{-\beta\omega/2})}{\frac{1}{2}\omega} \quad (3.3.22)$$

where

$$S_O(\omega) = \frac{1}{2\pi N} \int_{-\infty}^{\infty} dt e^{-i\omega t} \langle M, M(t) \rangle$$

and

$$\int_{-\infty}^{\infty} d\omega S_O(\omega) = Y \quad .$$

The integrand in (3.3.21) has a maximum at $\lambda=0$, from which it follows that

$$\chi_T \leq Y \quad .$$

The second inequality follows using the convexity of

$$\frac{1 - e^{-x}}{x} \quad .$$

3.4 Fourth order fluctuation

The gap parameter was defined in (2.1.7). To calculate the gap parameter one has to calculate the critical index, γ_2 , of

$$F_2 = \frac{1}{Nm^4 \beta^4} \left. \frac{\partial^4}{\partial H_x^4} \log Z \right|_{H_x=0}, \quad (3.4.1)$$

The gap parameter Δ is then

$$\Delta = \frac{1}{2} (\gamma_2 - \gamma) \quad . \quad (3.4.2)$$

We note that upon differentiation (3.4.1) becomes

$$Nm^4 \beta^4 F_2 = \frac{\left. \frac{\partial^4}{\partial H_x^4} \text{Tr } e^{\beta(JP + M_x H_x)} \right|_{H_x=0}}{\text{Tr } e^{KP}} - 3 \left[\frac{\left. \frac{\partial^2}{\partial H_x^2} \text{Tr } e^{\beta(JP + M_x H_x)} \right|_{H_x=0}}{\text{Tr } e^{KP}} \right]^2 \quad (3.4.3)$$

To obtain (3.4.3) we made use of the fact that above T_c and at $H=0$ odd derivatives of Z vanish.

When the magnetization and Hamiltonian commute, (3.4.3) reduces to

$$\beta^4 \{ \langle M_x^4 \rangle - 2 \langle M_x^2 \rangle^2 \} \quad . \quad (3.4.4)$$

The second term in (3.4.3) is proportional to $(\chi_{||}^0)^2$. In Section 3.3B we saw that $\chi_{||}^0$ behaves like $\langle M_x^2 \rangle$ as $T \rightarrow T_c$. The first term in (3.4.3) is

$$\frac{\beta^4 m^4}{Z} \sum_n \frac{K^{n-4}}{(n-4)!} \text{Tr} \sum_{ijkl} P^j_{M_x} P^k_{M_x} P^l_{M_x} P^i_{M_x} . \quad (3.4.5)$$

The difference between (3.4.5) and $\langle M_x^4 \rangle$ is in terms of commutators of M and \mathcal{H} . By physical reasoning this difference is expected not to diverge, or diverge less than the same quantity with the commutator brackets removed.

As mentioned in Section 3.3B commutators have the effect of removing two spins from each trace, thereby reducing the divergence.

A more rigorous treatment means going to the fourth order in the field in the expression for the density matrix

$$\begin{aligned} & \left(\frac{-1}{i\hbar}\right)^4 \int_{-\infty}^t \int_{-\infty}^{t_1} \int_{-\infty}^{t_2} \int_{-\infty}^{t_3} e^{\frac{it\mathcal{H}}{\hbar}} [M_x(t_1), [M_x(t_2), [M_x(t_3), [M_x(t_4), \rho]]]] \\ & \times e^{\frac{-it\mathcal{H}}{\hbar}} dt_1 dt_2 dt_3 dt_4 , \end{aligned} \quad (3.4.6)$$

which is rather complicated.

We treat therefore

$$Y_2 \approx \frac{16}{m} \frac{1}{N} \left\{ \frac{3}{2} \langle M_x^2 \rangle^2 - \frac{1}{2} \langle M_x^4 \rangle \right\} . \quad (3.4.7)$$

The definition (3.4.7) has the advantage of yielding a positive series starting with unity.

The term $\langle M_x^4 \rangle$ is a sum over four indices

$$\langle M_x^4 \rangle = \frac{m^4}{16} \sum_{i,j,k,\ell} \langle \sigma_i^x \sigma_j^x \sigma_k^x \sigma_\ell^x \rangle .$$

We consider all possible cases for the summand and list them below:

$$(i) \quad \langle \sigma_a^{x4} \rangle$$

$$(ii) \quad \langle \sigma_a^{x2} \sigma_b^{x2} \rangle , \quad a \neq b$$

$$(iii) \quad \langle \sigma_a^{x3} \sigma_b^x \rangle , \quad a \neq b$$

$$(iv) \quad \langle \sigma_a^{x2} \sigma_b^x \sigma_c^x \rangle , \quad a \neq b \neq c$$

$$(v) \quad \langle \sigma_a^x \sigma_b^x \sigma_c^x \sigma_d^x \rangle , \quad a \neq b \neq c \neq d$$

Summing (i) - (v) separately, we obtain

$$\begin{aligned} \langle M_x^4 \rangle = \frac{m^4}{16} \{ & N + 3N(N-1) + 8N \sum_{j \neq 0} \langle \sigma_0^x \sigma_j^x \rangle \\ & + 12N(N-2) \sum_{j \neq 0} \langle \sigma_0^x \sigma_j^x \rangle + \sum_{i \neq j \neq k \neq \ell} \langle \sigma_i^x \sigma_j^x \sigma_k^x \sigma_\ell^x \rangle \} \end{aligned} \quad (3.4.8)$$

$$\begin{aligned}
\langle M_x^2 \rangle &= \frac{m_N^2}{4} (1 + 2 \sum_{j \neq 0} \langle \sigma_0^x \sigma_j^x \rangle) \\
3 \langle M_x^2 \rangle^2 &= \frac{m_N^4}{16} (3 + 12 \sum_{j \neq 0} \langle \sigma_0^x \sigma_j^x \rangle + \\
&\quad + 12 \sum_{\substack{j \neq 0 \\ k \neq 0}} \langle \sigma_0^x \sigma_j^x \rangle \langle \sigma_0^x \sigma_k^x \rangle)
\end{aligned} \tag{3.4.9}$$

We see that the parts of the second and fourth terms in (3.4.7) which were proportional to N^2 are cancelled by the first two in (3.4.9). The third term of (3.4.9) subtracted from $\sum_{i \neq j \neq k \neq \ell} \langle \sigma_i^x \sigma_j^x \sigma_k^x \sigma_\ell^x \rangle$ leaves the part which is linear in N .

When $[M, \mathcal{H}] = 0$ there is no need to check on linearity in N , as F_2 is a derivative of the free energy which we know to be linear in N .

We can define a generalized exponential e_g^{A+B} in which factors B will always be taken to the front

$$e_g^{A+B} = e^A \cdot e^B.$$

Using this definition of order we see that

$$Y_2 = \frac{16}{Nm^4 \beta^4} \frac{\partial^4}{\partial H_x^4} \log_g \text{Tr} e_g^{-\beta(\mathcal{H} - M_x H_x)} \Big|_{H_x=0}$$

Kubo (1962) discussed these generalized exponentials and the calculation of generalized cumulants and moments with them. The relation between cumulants and moments

remains unchanged as long as we apply the generalized order consistently. Therefore

$$\log_g \text{Tr}_g e^{-\beta (\mathcal{H} - M_x H_x)}$$

is equal to the cumulant expansion series which is linear in N .

We see from (3.4.8) and (3.4.9) that

$$Y_2 = \frac{16m^4}{N} \left\{ 1 + 8 \sum_{j \neq 0} \langle \sigma_0^x \sigma_j^x \rangle - Q \right\} \quad (3.4.11)$$

where

$$Q = -\frac{1}{2} \sum_{j \neq j \neq k \neq \ell} \frac{K^n}{n!} \text{Tr}^* P^n \sigma_i^x \sigma_j^x \sigma_k^x \sigma_\ell^x.$$

The second term in (3.4.11) is four times the fluctuation series apart from the initial unity. We calculate the third term. We generate the graphs by opening two bonds in the partition function graphs, under the restriction that the four vertices be distinct.

We note that each graph $P^n a_a^\dagger a_b a_c^\dagger a_d$ appears 24 times in the sum as each index i, j, k, ℓ can take any value a, b, c , or d . On the other hand we are counting each term twice in the graphical expansion because the dotted arrow associated with $a_a^\dagger a_b a_c^\dagger a_d$ can be either

$a \rightarrow b, c \rightarrow d$ or $a \rightarrow d, c \rightarrow b$, for example



Therefore

$$Q = -6 \sum_n \frac{K^n}{n!} \sum_{g'_{n+2}} \frac{h(g'_{n+2}) V(g'_{n+2}) (g_n, L)}{2^S} . \quad (3.4.12)$$

The vertical weight is calculated with the dotted arrows fixed at the last two levels (since the four vertices are distinct their order is immaterial).

For elimination of equivalent graphs and calculation of horizontal weights the matrix elements $M(a,b)$, $M(c,d)$ are preceded by a minus sign. $N(a,b)$, $N(b,a)$, $N(c,d)$ and $N(d,c)$ are zero if the vertices $a-b$ and $c-d$ respectively are not connected by other arrows.

CHAPTER IV

THE SERIES FOR THE STATIC QUANTITIES AND
THEIR ANALYSIS4.1 Series analysis

The functions we deal with are assumed to have near T_c an asymptotic behaviour of the form:

$$f(x) \sim A(x_c - x)^{-(g+1)} . \quad (4.1.1)$$

We have a finite number of terms in a series expansion of $f(x)$ in powers of x . Two basic methods have been developed to obtain estimates of x_c and g from this truncated expansion

$$f(x) = \sum_{n=0}^N a_n x^n + R_N . \quad (4.1.2)$$

(a) The ratio method

This method was first used by Domb and Sykes (1957). Had (4.1.1) been an exact equality, using the binomial expansion we would see that the ratios r_n defined by

$$r_n = a_n / a_{n-1} \quad (4.1.3)$$

would be equal to

$$r_n = \frac{1}{x_c} \left(1 + \frac{g}{n}\right) . \quad (4.1.4)$$

If the series (4.1.2) converges, its radius of convergence in the complex x plane is the distance from the origin to the first singularity. The limit of $|r_n^{-1}|$, if it exists, is the radius of convergence. For series with all positive terms the first singularity must lie on the positive real axis and can be expected to be the physical singularity we are seeking; therefore we expect the ratios to approach x_c^{-1} as $n \rightarrow \infty$. Since we expect (4.1.1) to hold near x_c the relation (4.1.4) will obtain in the limit $n \rightarrow \infty$. A plot of r_n against $1/n$ will therefore approach a straight line with intercept $1/x_c$ at $\frac{1}{n}=0$ and a limiting slope of g .

Various modifications and refinements of the ratio method have been introduced. Neville tables have been used by Baker, Eve, Gilbert and Rushbrooke (1968) which are successive linear extrapolations. The quantities l_n , q_n etc. are constructed $l_n = n r_n - (n-1) r_{n-1}$, $q_n = \frac{1}{2} [n l_n - (n-2) l_{n-1}]$, $c_n = \frac{1}{3} [n q_n - (n-3) q_{n-1}]$, etc., and as long as they are monotonic they provide improved estimates of x_c^{-1} . Once the critical point is evaluated, better estimates for the exponent can be obtained by the relation (Domb and Sykes (1957))

$$n [x_c (a_n/a_{n-1}) - 1] \rightarrow g \quad \text{as } n \rightarrow \infty.$$

If this sequence is monotonic a Neville table can be

constructed for it too. Given an estimate g' of g one can obtain a better estimate for x_c from the sequence

$$r_n(n/(n + g')). \quad (4.1.5)$$

Alternate ratios are often used to smooth out oscillations for loose packed lattices. Fisher (1962) suggested the sequences

$$r'_n = \frac{1}{2} \{ (n+\epsilon) r_n - (n+\epsilon-2) r_{n-2} \} \quad (4.1.6)$$

for a choice of small ϵ .

In general the series will have non physical singularities too; the ratio sequence will converge more rapidly when they are far from the circle of convergence.

For series of alternating sign the dominant singularity will be on the negative real axis and for this and for the general non positive series the ratio method will not yield the physical singularity. On the other hand a pair of singularities close to the real axis could show as a consistent limit while the series is well behaved on the real axis and has no physical singularity, but for three dimensional cases we expect a physical singularity so this problem occurs only in the two dimensional case. Transformations of variable, usually bilinear, have been used to bring the physical singularity in to be the

closest to the origin (Wortis (1969), Guttman, Thompson (1969)), with noticeable success.

(b) Pade approximants

The (N,M) Pade approximant to a power series $F(x)$ is the ratio of two polynomials of degree N and M , whose $N+M+1$ coefficients are determined uniquely by the requirement that the coefficients of the power expansion of the ratio be equal to those of $F(x)$ up to the $N+M$ order.

$$[N,M] = \frac{P(x)}{Q(x)} = \frac{p_0 + p_1 x + \dots + p_N x^N}{1 + q_1 x + \dots + q_M x^M} \quad (4.1.7)$$

First introduced to critical phenomena by Baker (1961) it has proved an invaluable tool for the analysis of series and especially for non positive series where the ratio method does not apply.

In a series of papers Baker has shown that for certain classes of functions the $[M,M]$ and $[M-1,M]$ approximants form upper and lower bounds to $F(x)$ for real positive x and that subsequences of the diagonal approximants converge to $F(z)$ everywhere in the complex plane as $M \rightarrow \infty$.

The zeroes of the denominator are calculated and the singularity of $F(x)$ shows up consistently, usually as the smallest positive zero. Obviously the approximation is better if the function's singularity is a simple

pole; therefore it is customary to calculate Pade approximants to

$$(i) \quad \frac{\partial}{\partial z} \log F(z),$$

or if an estimate of γ is known to

$$(ii) \quad [F(z)]^{1/\gamma}, \quad \text{and}$$

$$(iii) \quad \left[\frac{\partial}{\partial z} F(z) \right]^{1/\gamma+1}.$$

We have assumed

$$F(z) = (z_c - z)^{-\gamma} G(z) \quad (4.1.8)$$

with $G(z)$ a non singular function at z_c . It follows immediately that

$$(i) \quad \frac{z}{dz} \log F(z) = - \frac{\gamma}{z-z_c} + \frac{d}{dz} \log G(z) \quad (4.1.9)$$

So if $G(z)$ has no zeroes near z_c that pole should appear with good convergence in the Pade tables.

$$(ii) \quad (F(z))^{1/\gamma+\epsilon} = (z_c - z)^{-1-\epsilon} (G(z))^{1/\gamma+\epsilon} \quad (4.1.10)$$

Here the convergence is better as the zeroes of $G(z)$ will not interfere. (iii) with an approximate γ is:

$$\begin{aligned} \left(\frac{\partial F(z)}{\partial z} \right)^{1/\gamma+1+\epsilon} &= (-\gamma G(z))^{1/\gamma+1+\epsilon} (z_c - z)^{-1-\epsilon} + \\ &+ \left(\frac{\partial G}{\partial z} \right)^{1/\gamma+1+\epsilon} (z_c - z)^{-\gamma/\gamma+1} \end{aligned} \quad (4.1.11)$$

Here the convergence will be slower than in (4.1.10) because of the second term but still better than (4.1.9). When the estimates for γ are not close to the real value one expects a shift in the estimate for z_c from Pade approximants to (4.1.10) and (4.1.11) since

$$(z_c - z)^{-1+\varepsilon} g(z) = (z_c)^{-1+\varepsilon} g(z) \left[1 + (\varepsilon-1) \frac{z}{z_c} + O\left(\frac{z}{z_c}\right)^2 \right]. \quad (4.1.12)$$

So to a first approximation the pole of this junction is at

$$y_c \approx \frac{z_c}{1-\varepsilon} \approx z_c (1 + \varepsilon)$$

varying the value of ε and plotting y_c against ε . We have a straight line, for ε small enough (in practice up to rather large values), whose intercept with the correct value of z_c will give a very good estimate of the exponent.

Another way of estimating the exponents is to calculate the Pade approximant at z_c to

$$(z-z_c) \frac{d}{dz} \ln F(z) = \gamma + O(z-z_c) \quad . \quad (4.1.13)$$

A method which is less sensitive to the choice of z_c was introduced by Baker, Eve, Gilbert and Rushbrooke (1968); the Pade approximant at z_c to

$$\frac{\frac{d}{dz} \ln \frac{dF(z)}{dz}}{\frac{d}{dz} \ln F(z)} \quad (4.1.14)$$

gives an estimate for $\frac{\gamma+1}{\gamma}$. Since (4.1.14) is $\frac{\gamma+1}{\gamma}$ at all z not far from z_c , this is insensitive to the choice of z_c . A plot of the Pade approximants to (4.1.14) gives a straight line in a γ against z_c plot. Its intercept with the straight line obtained from (4.1.10) yields a simultaneous estimate of γ and z_c .

4.2 The static series and their analysis

Using the graphical expansions outlined in Chapter III, we have obtained the initial parallel susceptibility and fourth order fluctuation high temperature series for the face-centered cubic and triangular lattices. This choice of lattices was made because the close packed lattices are known to yield better behaved high temperature series.

For the Heisenberg model it is also desirable to study the staggered susceptibility on loose packed lattices in order to locate the Néel point, T_N . However there is no reason to do these calculations on the loose packed lattices since for the XY model the staggered parallel susceptibility is equal to the parallel ferromagnetic susceptibility.

The series obtained are:

for the face centered cubic lattice,

$$\bar{\chi}_\parallel^O = 1 + 6K + 32K^2 + 161.5K^3 + 792.3K^4 + 3823.924K^5 + \\ + 18262.124K^6 + 86567.462K^7 + 402842.0556K^8 + \dots$$

$$Y_2 = 1 + 24K + 339K^2 + 3656K^3 + 33176.25K^4 + 268835K^5 + \\ + 2010189.817K^6 + 13768318.21K^7 + \dots$$

for the triangular lattice,

$$\bar{\chi}_\parallel^O = 1 + 3K + 7K^2 + 13.75K^3 + 25.25K^4 + 41.0125K^5 + \\ + 67.6702K^6 + 109,50186K^7 + 170.5454K^8 + \dots$$

$$Y_2 = 1 + 12K + 79.5K^2 + 379K^3 + 1432.875K^4 + 4621.5K^5 + 13336.25K^6 + 35115.962K^7 + \dots$$

(a) Comparison of susceptibility and fluctuation

In view of the asymptotic equivalence of the fluctuation and reduced susceptibility series we wish to compare the ratios of the coefficients.

TABLE 4.1

Ratio of fluctuation over susceptibility coefficients

n	triangular lattice	f.c.c. lattice
1	1	1
2	1.071428571	1.03125
3	1.05454545	1.01541799
4	.965346534	1.00435441
5	.971959768	1.00149088
6	.985478098	1.00111803
7	1.011481242	1.00089795
8	1.026660712	1.01355544

In the case of the f.c.c. lattice the ratios are approaching unity ever more closely as n increases

(except for the last term). The last term breaks the trend both here and in the ratio plot and we suspect a numerical error in the susceptibility series might be the cause.

In the triangular lattice case the ratios of coefficients of susceptibility over fluctuation are more erratic. From the result of Falk and Bruch $\chi \leq Y$ at all temperatures therefore the ratios should be greater or equal to unity as n increases and this seems to be the trend.

(b) Analysis of the f.c.c. susceptibility series

Figure 4.1 is the ratio plot for the f.c.c. susceptibility series. We see that the ratio plot is linear (except for the last term) and we can estimate

$$K_C^{-1} = 0.2206 \pm .001 \quad \gamma = 1.33 \pm 0.03 .$$

The above results agree well with the ratio analysis of the fluctuation (Betts, Elliott and Lee (1970)) who find $K_C^{-1} = 0.2206 \pm .0010$ and $\gamma = 1.335 \pm 0.02$.

We apply the modified ratio method (Section 4.1) to the f.c.c. susceptibility series.

TABLE 4.2

Ratios of f.c.c. susceptibility with $K'_C=0.24$, $\gamma'=1.33$

n	r_n	$n(K'_C r_n - 1) + 1$	$n a_n / (n + \gamma' - 1) a_{n-1}$
1	6	1.32600	4.51128
2	5.33333	1.35733	4.57798
3	5.046875	1.36058	4.54673
4	4.905882	1.33680	4.56023
5	4.826358	1.33313	4.52754
6	4.775754	1.33265	4.52678
7	4.740273	1.33320	4.52686

Fig. 4.2 exhibits:

- (1) The physical pole of $(\chi_{||}^O(K))^{1/\gamma}$ as a function of γ .
- (2) The estimate of γ from $(K-K'_C) \frac{\partial}{\partial K} \log \chi_{||}^O(K) \Big|_{K=K'_C}$ against K'_C .

This is given by the vertical bars. The bars shrink as approach K'_C .

- (3) The arrow at the left hand side shows the spread in γ calculated from Pade approximants to:

$$\frac{\partial}{\partial K} \log \frac{\partial}{\partial K} \chi_{||}^O(K) / \frac{\partial}{\partial K} \log \chi_{||}^O(K) .$$

The approximants to $(\frac{\partial \chi}{\partial K})^{1/\gamma+1}$ gave a straight line when plotting the pole against γ . The line had the same slope as (1) only with larger spread. We also calculated Padé approximants to $\frac{\partial}{\partial K} \log \chi(K)$ listed in Table 4.3, and to $(\frac{\partial^2}{\partial K^2} \chi/\chi)^{1/2}$ listed in Table 4.8. The results are all consistent.

TABLE 4.3

Estimates of K_c for the f.c.c. lattice from
Padé approximants to $\frac{\partial}{\partial K} \log \chi_{II}(K)$.

N	M = 2	M = 3	M = 4	M = 5	M = 6	M = 7
0	.2255	.2233	.2213	.2207	.2206	.2230± .005i
1	.2236	.2095	.2203	.2206	.2207	
2	.2283	.2105	.2208	.2203		
3	.2201	.2208	.2205			
4	.2201	.2201				
5	.2206					

TABLE 4.4

Estimates of K_c for the f.c.c. lattice
from Pade approximants to $(\chi_{||}(K))^{4/5}$

N	M = 2	M = 3	M = 4	M = 5	M = 6	M = 7
1	.20017	.21911	.21887	.21910	.21973	.21882
2	.21895	.21886	.21898	.21820	.21909	
3	.21886	.21892	.2147± 0.005i	.21897		
4	.21911	.21722	.21892			
5	.21955	.21911				
0	.21865					

TABLE 4.5

Estimates of K_c for the f.c.c. lattice
from Pade approximants to $(\chi_{||}(K))^{3/4}$

N	M = 2	M = 3	M = 4	M = 5	M = 6	M = 7	M = 8
0		.2206	.2212	.2212	.2211	.2210	.2271
1	.2196	.2216	.2212	.2212	.2208	.2211	
2	.2215	.2212	.2210	.2210	.2212		
3	.2212	.2210	.2210	.2210			
4	.2212	.2210	.2210				
5	.2207	.2212					
6	.2211						

TABLE 4.6

Estimates of K_C for the f.c.c. lattice
from Pade approximants to $(\chi_{II}(K))^{10/13}$

N	M = 2	M = 3	M = 4	M = 5	M = 6	M = 7
1				.22022	.22039	.22020
2	.22045	.22019	.22021	.22023	.22022	
3	.22021	.22021	.22018	.22021		
4	.22021	.22021	.22021			
5	.22033	.22021				
6	.22017					

TABLE 4.7

Estimates of γ for the f.c.c. lattice from
Pade approximants to $(K-K_C) \frac{\partial}{\partial K} \log \chi_{II}(K) \Big|_{K=K_C}$

Pade	$K_C=.2200$	$K_C=.2202$	$K_C=.2204$	$K_C=.2206$	$K_C=.2208$	$K_C=.221$	$K_C=.222$
(1,3)	1.293			1.31391		1.327	1.356
(2,2)	1.294	1.3014	1.30845	1.31535	1.32204	1.329	1.358
(2,3)	1.291	1.2995	1.30772	1.31579	1.32372	1.331	1.368
(3,2)	1.291	1.2994	1.30770	1.31579	1.32363	1.331	1.365
(1,4)	1.291	1.2994	1.30764	1.31521	1.32213	1.328	1.349
(2,4)	1.294	1.3001	1.30721	1.31507	1.32371	1.333	1.393
(4,2)	1.291	1.2998	1.30666	1.31497	1.32371	1.333	1.386
(3,3)	1.300	1.3037	1.30879	1.31547	1.32371	1.334	1.408
(1,5)	1.232	1.2706	1.29872	1.31495	1.31862	1.310	1.119
(3,4)	1.291	1.2995	1.30772	1.31580	1.32372	1.331	1.368
(4,3)	1.291	1.2994	1.30771	1.31579	1.32363	1.331	1.365
(2,5)	1.291	1.2995	1.30765	1.31521	1.32213	1.328	1.350
(5,2)	1.291	1.2993	1.30731	1.31474	1.32163	1.328	1.351
(1,6)	1.295	1.3014	1.30828	1.31525	1.32230	1.329	

TABLE 4.8

Estimates of K_c for the f.c.c. lattice from
 Pade approximants to $\left[\frac{\partial^2}{\partial K^2} \chi_{||}(K) / \chi_{||}(K)\right]^{\frac{1}{2}}$.

N	M = 2	M = 3	M = 4	M = 5	M = 6
0		.2217	.2206	.2206	.2145±.06i
1		.2197	.2206	.2206	
2	.2193	.2207	.2197		
3	.2206	.2193			
4	.2206				

(c) Analysis of the fourth order fluctuation series
 for the f.c.c. lattice

The ratio plot is given by the solid line in
 Fig. 4.3. The modified ratios using the known value
 of K_c are tabulated below and plotted in Fig. 4.3
 (the dotted line).

TABLE 4.9

Ratios of $Y_2(K)$ for f.c.c. lattice, $K'_c = .221$

n	r_n	$(K'_c r_n - 1) n + 1$
2	14.125	5.243
3	10.785	5.150
4	9.074	5.022
5	8.103	4.954
6	7.477	4.915

We estimate then $\gamma_2 = 4.64 \pm 0.10$ if we ignore the last K' term.

The Pade approximants to the fourth order fluctuation are not well behaved. A few of the approximants are tabulated below. There was no point in plotting the equivalent of Fig. 4.1 as the spread in estimates from different Pades was large, and the K_c against γ plot from the estimates to $[Y_2(K)]^{1/\gamma}$ almost vertical.

Tables 4.11 and 4.12 are Pade approximants to $\frac{\partial}{\partial K^x} \log Y_2(K^x)$ and $(K^x - K_c^x) \frac{\partial}{\partial K^x} \log Y_2(K^x)$ with $K = \frac{0.5K^x}{1-0.5K^x}$ which seemed somewhat better behaved. All the above were considered with the estimate from the modified ratio method in table 4.9.

TABLE 4.10

Estimates of K_c for f.c.c. lattice from
Pade approximants to $[Y_2(K)]^{1/4.65}$

N	M = 1	M = 2	M = 3	M = 4	M = 5	M = 6
1	.2125	*	*	.2376	.2113	.1919
2	.2269	*	*	.2249	.2317	
3	.2423	.2302	.2231	.2387		
4	.1908	.2127	.2281			
5	.2318	.1853				
6	*					

* The physical pole did not appear at all or was off the real axis.

TABLE 4.11

Estimates of K_c for the f.c.c. lattice from

Pade approximants to $\frac{\partial}{\partial K^X} \log Y_2(K^X)$ where $K = \frac{\frac{1}{2}K^X}{1 - \frac{1}{2}K^X}$

N	M = 2	M = 3	M = 4
2	.4011	.3634	.3702
3	.3159	.3697	
4	*		

TABLE 4.12

Estimates of γ_2 for the f.c.c. lattice from

Pade approximants to $\left[(K^X - K_c^X) \frac{\partial}{\partial K^X} \log Y_2(K^X) \right] \Big|_{K^X = .356}$
 where $K = \frac{\frac{1}{2}K^X}{1 - \frac{1}{2}K^X}$

N	M = 1	M = 2	M = 3	M = 4	M = 5
1	4.5188	6.5715	4.0173	4.6461	4.6300
2	5.2479	4.3449	4.2448	4.6307	
3	3.9788	4.2378	4.3037		
4	4.7663	4.7461			
5	4.7428				

TABLE 4.13

Estimates of K_c for the f.c.c. lattice
from Pade approximants to $\frac{\partial}{\partial K} \log Y_2(K)$

N	M = 2	M = 3	M = 4	M = 5
1	*	.2754	.1938	.2272
2	.2694	.2440	.2474	
3	.164	.2473		
4	*			

TABLE 4.14

Estimates of γ_2 for the f.c.c. lattice from
Pade approximants to $(K-K_c) \frac{\partial}{\partial K} \log Y_2(K) \Big|_{K_c=.221}$

N	M = 2	M = 3	M = 4	M = 5
1	6.367	4.151	5.477	18.436
2	5.473	4.502	4.240	
3	4.437	2.969		
4	4.075			

Fig. 4.1

The ratio plot of the initial parallel
susceptibility for the f.c.c. lattice

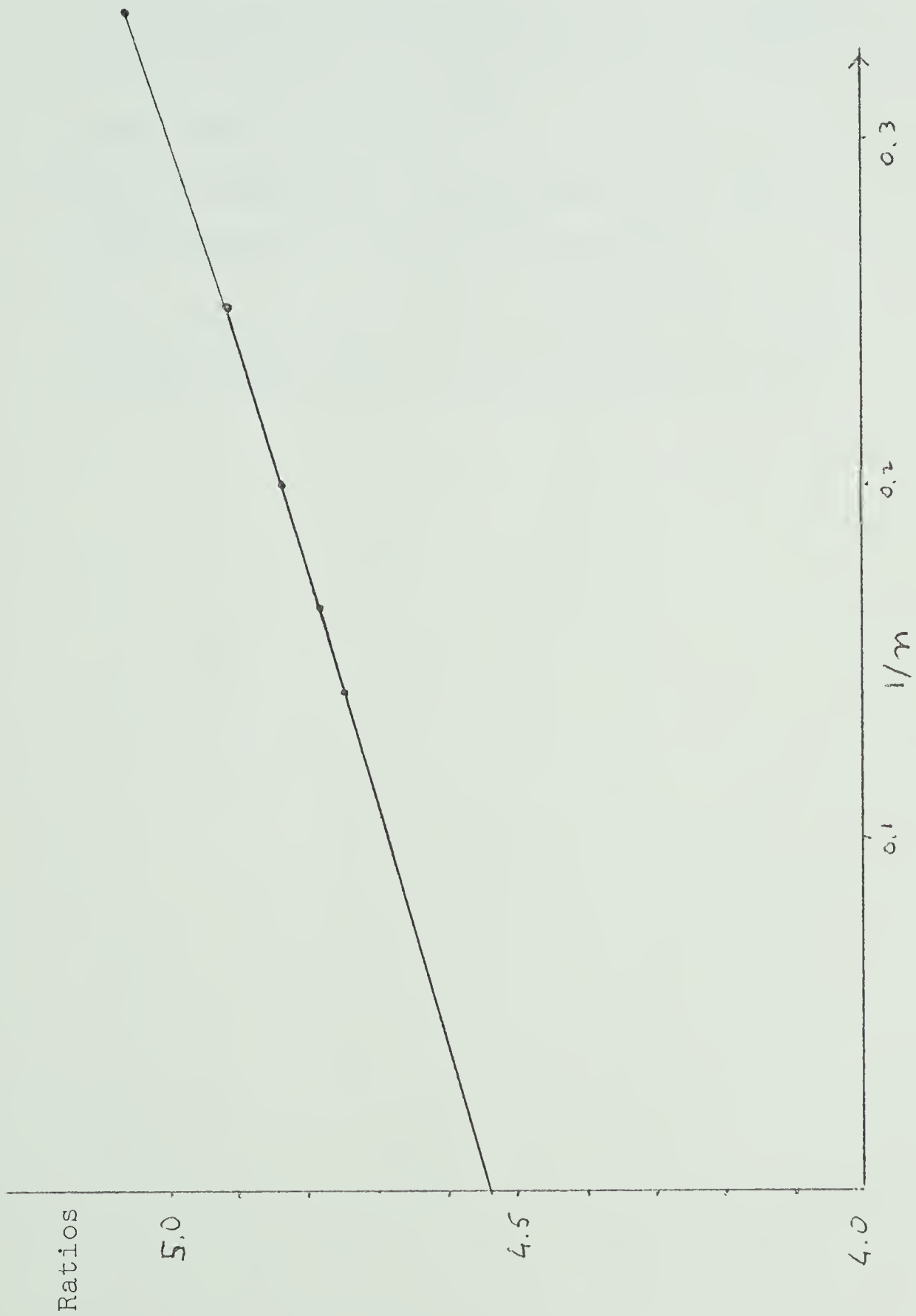


Fig. 4.2

Estimate of K_c and γ for the f.c.c. lattice.

The solid line is the estimate of K_c from approximants to $[\chi(K)]^{1/\gamma}$.

The vertical bars are estimate of γ from approximants to $(K-K_c) \frac{\partial}{\partial K} \log \chi \Big|_{K=K_c}$.

The arrow on the left is estimate of γ from $\frac{\partial}{\partial K} \log \chi / \frac{\partial}{\partial K} \log \chi$ at various values of K .



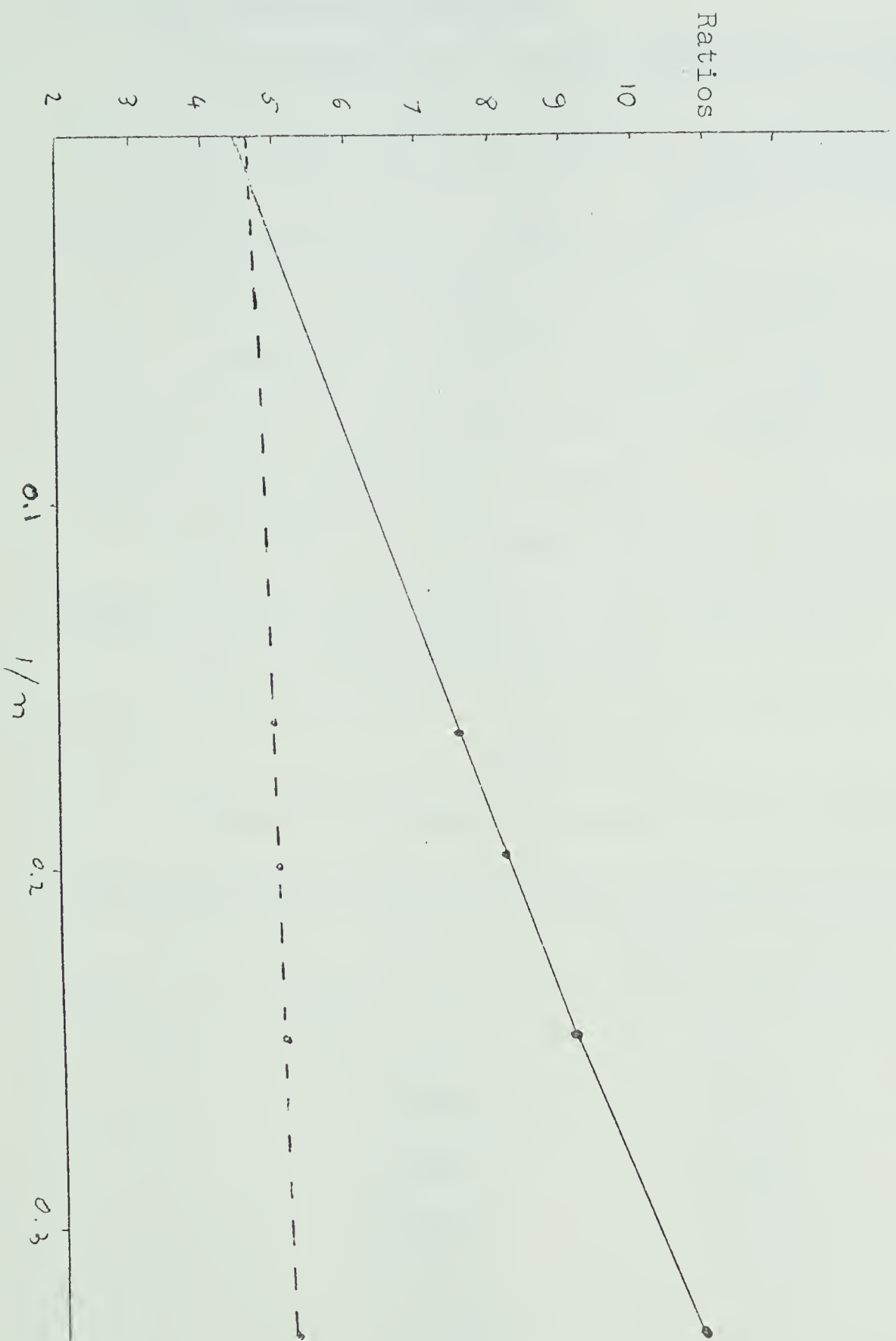
Fig. 4.3

The ratio plot for the fourth order fluctuation of the f.c.c. lattice.

The solid line is the ratio plot of the series.

The dotted line is the plot of the estimates of γ

$$\gamma_n = (K_c r_n - 1) n - 1.$$



(d) The triangular lattice

The evidence for a phase transition in the XY model on the triangular lattice is not conclusive. We present our results below; a discussion is presented in the next section of conclusions.

In Fig. 4.4 the plot of the physical singularity from approximants to $(\chi^0)^{1/\gamma}$ against γ is given. On the same Fig. 4.4 the estimate of γ from the approximants to $\frac{\partial}{\partial K} \log \frac{\partial}{\partial K} \chi / \frac{\partial}{\partial K} \log \chi$ is shown.

From Fig. 4.4 we could say that if there is a phase transition it occurs at $K_c = .65 \pm .07$.

Fig. 4.5 is the ratio plot.

TABLE 4.15

Ratios of triangular susceptibility series $K'_c = \frac{2}{3}$

n	r_n	$n(K'_c r_n - 1) + 1$
1	3	2
2	2.33333	2.1111
3	1.96429	1.9286
4	1.83636	1.8970
5	1.62426	1.4142
6	1.64929	1.0000
7	1.61817	1.5515

TABLE 4.16

Ratios of $\frac{\partial}{\partial K} \chi_n(K)$ for the triangular lattice,
 $K'_c = \frac{2}{3}$

n	r_n	$M (K'_c r_n - 1) + 1$
1	4.66667	3.1111
2	2.94642	2.9286
3	2.44847	2.8969
4	2.03032	2.4142
5	1.98000	2.600
6	1.88786	2.551

The Pade approximants to $\frac{\partial}{\partial K} \log \chi(K)$ are not converged at all.

TABLE 4.17

Estimates of K_c for the triangular lattice
 from Pade approximants to $\frac{\partial}{\partial K} \ln \chi_n(K)$

N	M = 1	M = 2	M = 3	M = 4	M = 5	M = 6	M = 7
0		*	.5924	*	.4814	*	.5717
1	.9524	.7502	*	*	.5969	*	
2	.5833	.8552	*	*	.6552		
3	*	*	.8871	.7888			
4	*	.3557	.5988				
5	.9998	*					

TABLE 4.18

Estimates of γ for the triangular lattice from
 Pade approximants to $\left[(K-K_C) \frac{\partial}{\partial K} \ln \chi(K) \right] \Big|_{K_C=2/3}$

N	M = 2	M = 3	M = 4	M = 5
1	1.8219	1.6724	1.3187	1.6427
2	2.1888	6.2717	1.4412	
3	3.6626	2.0076		
4	1.2286			

The analysis of the fourth order fluctuation is not too conclusive. The series itself seems to diverge at $K_C^{-1} = 1.2 \pm 0.2$ which could be interpreted both as $K_C = 1$ which is just the mathematical radius of convergence for a positive series or as a physical singularity at the lower limit.

TABLE 4.19

Ratios of the fourth order fluctuation on
 the triangular lattice, $K'_C = 2/3$

n	r_n	$n(K'_C r_n - 1) + 1$
1	12	6
2	6.625	7.8333
3	4.76730	7.53460
4	3.78067	7.08179
5	3.22533	6.75110
6	2.88570	6.54280

The ratio plots are given in Fig. 4.6.

TABLE 4.20

Estimates of γ_2 for the triangular lattice from
Pade approximants to $(K-K_c) \left. \frac{\partial}{\partial K} \log Y_2(K) \right|_{K_c=0.6667}$

N	M = 2	M = 3	M = 4
2	8.1894	5.0272	5.0512
3	4.1519	5.0500	
4	5.7267		

The value of $K_c = 0.6667$ was obtained by a better analysis of the fluctuation series (Betts - private communication).

Fig. 4.4

Ratio plot for the parallel initial susceptibility
on the triangular lattice

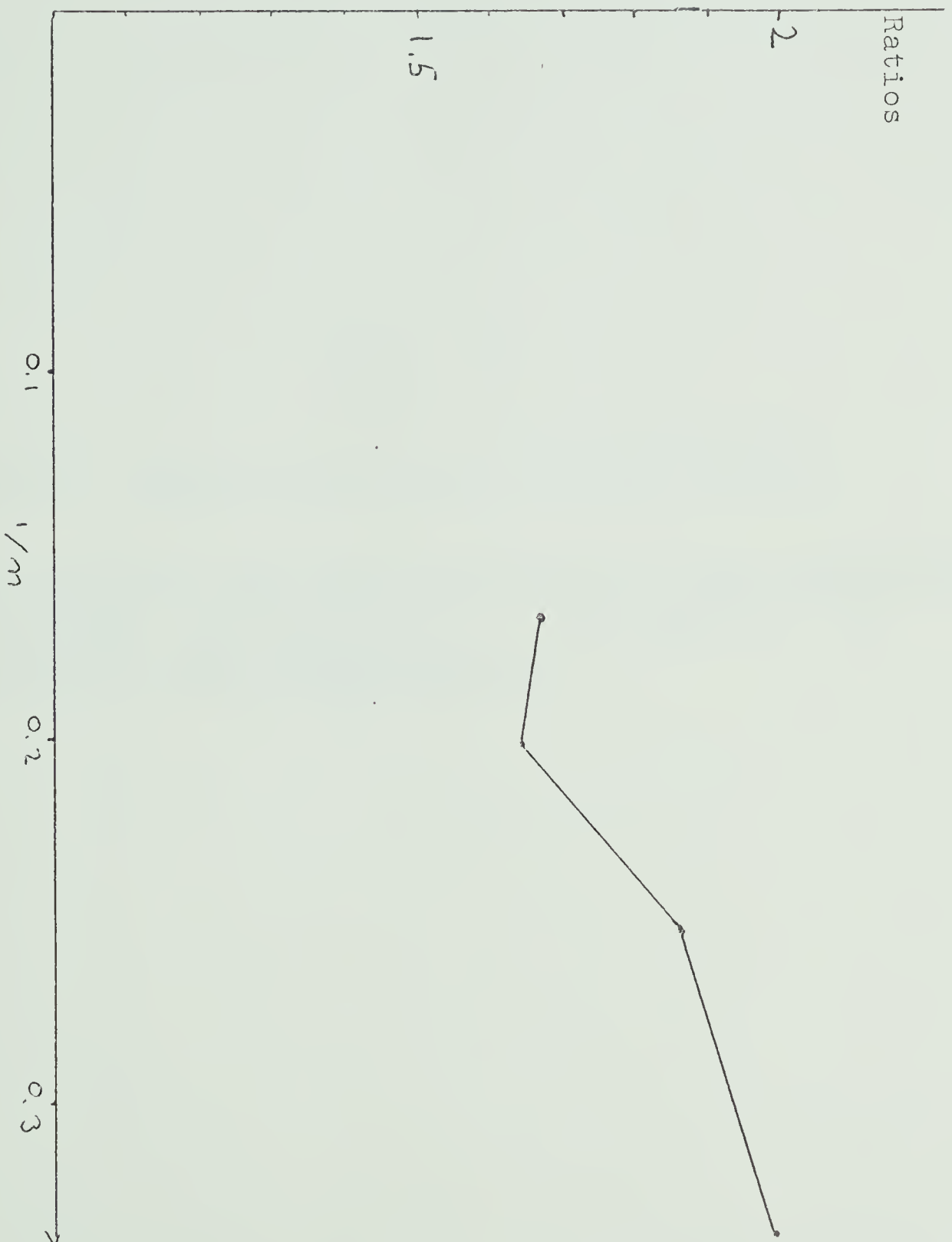


Fig. 4.5

Plot of $\chi(K)$ against K_c for the triangular lattice.

The horizontal bars show the spread of Padé approximants to $[\chi(K)]^{1/\nu}$. The vertical bar is the spread in χ as estimated from $\frac{\partial}{\partial K} \log \chi$ and $\frac{\partial}{\partial K} \log \chi$.

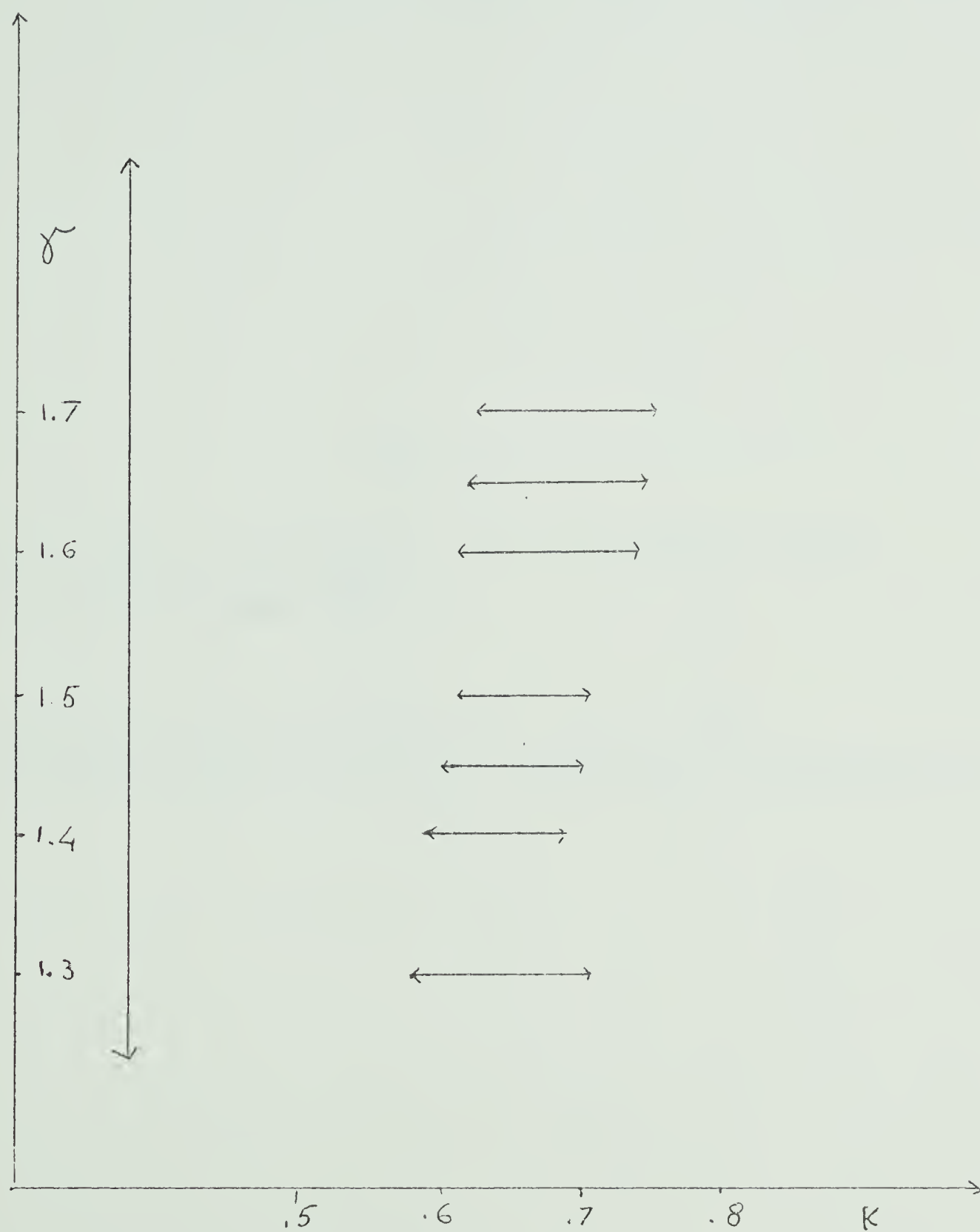
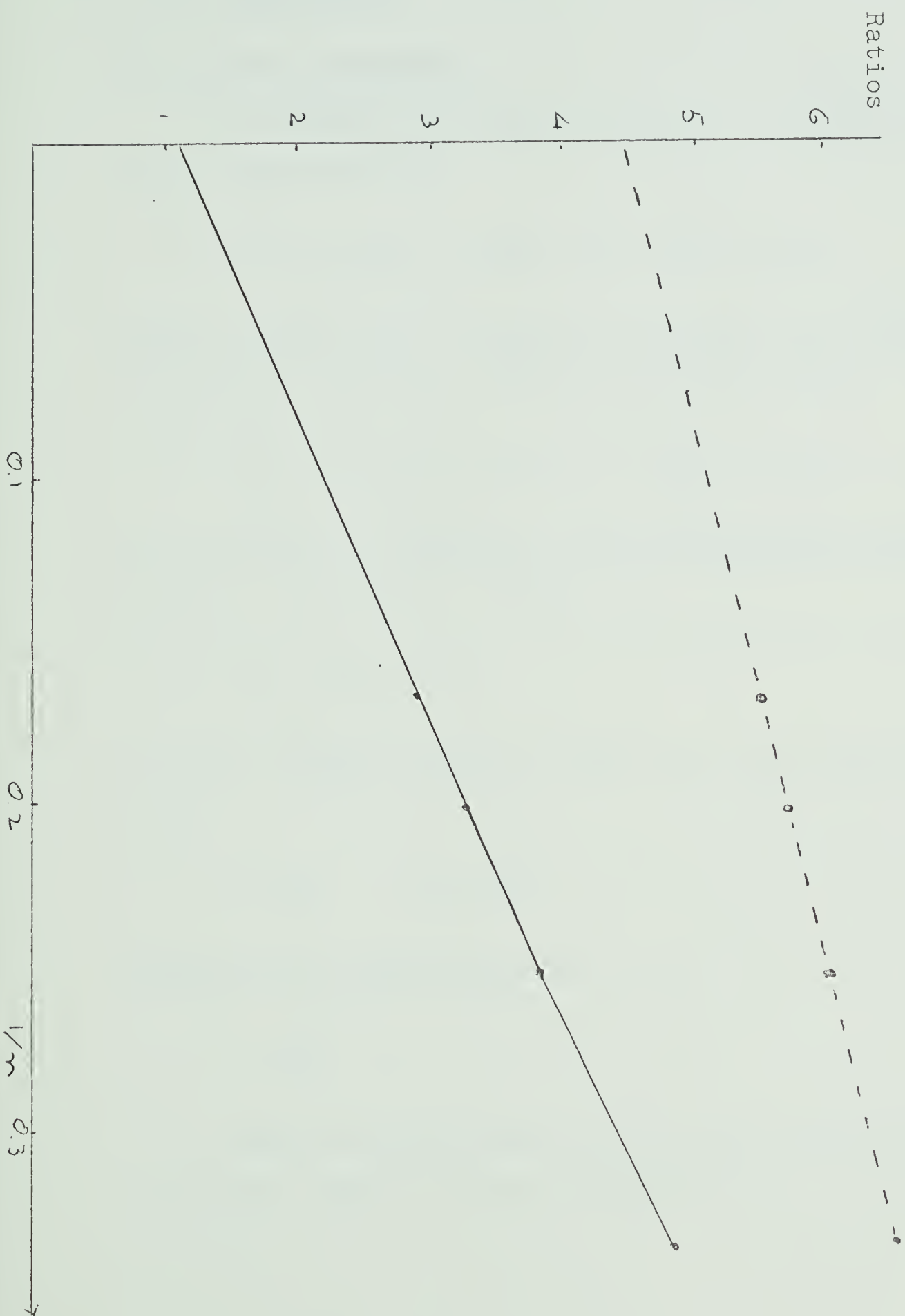


Fig. 4.6

Ratio plot of the fourth order fluctuation series
on the triangular lattice.

The solid line is the regular a_n/a_{n-1} plot.

The dotted line is the plot of $(\tilde{\gamma}-1)_n = (K_c r_n - 1)n$ against $1/n$.



4.3 Conclusion

(a) Three dimensions

We obtain from the susceptibility series on the f.c.c. lattice that

$$K_C = 0.2206 \pm 0.001 \quad \gamma = 1.315 \pm 0.04 ,$$

which agrees with the results obtained from the fluctuation series by Betts, Elliott and Lee (1970)

$$K_C = 0.2210 \pm 0.0005 \quad \gamma = 1.35 \pm 0.03 .$$

The analysis of the fourth order fluctuation yields (Ditzian and Betts (1970))

$$\gamma_2 = 4.64 \pm 0.1 .$$

Possible simple fractions consistent with our estimates are :

$$\gamma = \frac{4}{3} \quad \gamma_2 = \frac{14}{3} .$$

Therefore the gap parameter Δ is :

$$\Delta = \frac{5}{3}$$

Using scaling relations (Section 2.4) we obtain all other indices listed in Table 4.21.

TABLE 4.21

Index	Scaling relation	Prediction	Experiments for He
γ		4/3	
Δ		5/3	
α	$2 + \gamma - 2\Delta$	0	0.000 ± 0.003
β	$\Delta - \gamma$	1/3	0.333 ± 0.010
δ	$\Delta / (\Delta - \gamma)$	5	
ν	$(2\Delta - \gamma) / 3$	2/3	0.67 ± 0.04
η	$(4\Delta - 5\gamma) / (2\Delta - \gamma)$	0	

The experimental results are taken from the following experiments:

The specific heat exponent α has been measured by Ahlers (1967), who measured C_p to great accuracy taking into account gravitational field effects and vaporization. The temperature dependence of the superfluid density in He^4 near T_λ has been measured by Clow and Reppy (1966). The quantity measured is the angular momentum caused by a persistent current. The result is a value for β the exponent of the order parameter.

The value of ν has been measured by Henkel, Smith and Reppy (1969). The experiment consisted in measuring the angular momentum caused by a persistent current

in a film of known thickness D . The angular momentum is linear in D and becomes zero at D_0 . The value of $D_0 - D$ solid is the healing length ξ_s under the assumption that near the critical point only one coherence length exists

$$\xi_s \propto (T - T_c)^{-\nu}.$$

We see an excellent agreement of the XY model with experiments on He^4 .

Domb (19) and Domb and Hunter (1965) advanced on Taylor series approach which predicted that the exponent δ be an odd integer. Our result of $\delta = 5$ agrees and also supports the conjecture that $\delta = 5$ for all three dimensional models.

(b) Two dimensions

The Ising model in two dimensions has not only a phase transition but the only analytically demonstrated phase transition. The other models in two dimensions are in a much more dubious state.

Mermin and Wagner (1966) proved that spontaneous magnetization cannot occur in one and two dimensions for the isotropic Heisenberg model. They show, using the Bogoliubov inequality (1962), that in two dimensions

$$|M(H, T)| < B T^{-\frac{1}{2}} |\ln |H||^{\frac{1}{2}}$$

where B is a constant.

This proof, as we mentioned in 3.3B does not apply to the XY model.

Stanley and Kaplan (1966, 1967) have shown that zero magnetization at all $T > 0$ does not imply to phase transition. The correlations can be of a qualitative longer range below some T_c and the susceptibility may diverge even if $M = 0$.

Stanley and Kaplan (1966, 1967), Moore (1969) show for the classical Heisenberg model and plane model clear evidence of phase transition. Stanley (1969) has evidence of phase transition in two dimensions for the quantum Heisenberg model.

The suggested behaviour, Dymn (unpublished), Stanley (1967, 1969), if

$$\langle S_O S_R \rangle \propto R^{-\lambda T}$$

such that $M = 0$ yet $\chi = \infty$.

Our evidence is not really conclusive. It seems to indicate a divergence in the susceptibility at $K_c \sim 0.67 \pm 0.1$ with an exponent $\gamma = 1.5 \pm 0.2$ which agree with more accurate results of the fluctuation

(Betts - private communication) which seems to diverge at $K_c = .67 \pm .01$ with exponent $\gamma = 1.5 \pm 0.02$.

The fourth order fluctuation series by itself does not help in deciding whether this is a phase transition or not as its critical point is not well defined and it could be that the series converges up to $K = 1$. If there is a phase transition at $K_c = 0.667$ then from the modified ratio method $\gamma_2 = 5.4 \pm 0.5$. Possible ratios within the error are

$$\gamma = \frac{12}{8} \quad \gamma_2 = \frac{42}{8}$$

$$\text{then } \Delta = \frac{15}{8}, \quad \delta = 5, \quad \nu = \frac{9}{8}$$

$$\beta = \frac{3}{8}, \quad \eta = \frac{2}{3}, \quad \alpha = -\frac{2}{8}.$$

CHAPTER V

DYNAMICAL CRITICAL PHENOMENA

5.1 Introduction

The critical region is usually characterized by long range and long time fluctuations. Reaching equilibrium therefore takes a long time and it is natural to extend critical phenomena theories to cover dynamic properties.

In order to measure dynamical properties of the system one has to use probes such as neutron scattering or frequency dependent magnetic fields. While neutron scattering experiments depend on both transfer of energy $\hbar\omega$ and transfer of momentum $\hbar\mathbf{k}$, frequency dependent fields have infinite wavelength, ($\mathbf{k}=0$). For this reason frequency dependent measurement yields only results that do not probe into the spatial dependence of the correlations. Since our calculation did not include neutron scattering we will not discuss it.

In the Ising model correlations of the z components of the spins can only have space dependence as each σ_i^z commutes with the Hamiltonian and therefore is constant in time. Fisher and Burford (1967)) calculated the zz correlations, which give the elastic scattering of neutrons from an Ising model.

Allan and Betts (1968) using the Kubo linear response method calculated exactly the frequency dependent susceptibility $\chi^{xx}(\omega)$ and inelastic neutron scattering from a spin $\frac{1}{2}$ Ising model in 2 dimensions.

Essam and Garelick (1968) obtained the exact solution for all spins and in presence of parallel fields, their solution was based on solving the Green's function hierarchy of operations of motions without need for approximations other than linear response.

Allan and Betts (1968) noticed that the limit of $\chi^{xx}(\omega)$ as ω goes to zero was equal to $\chi_T^{xx}(0)$ in the case of the honeycomb lattice but not the square lattice. The Kubo "adiabatic", or quasistatic susceptibility is not equal to the isothermal susceptibility on the square lattice because the system is not ergodic in the absence of a field in the x direction.

As shown by Essm and Garelick (1968) $M = - \sum_r P_r E_r'(H)$ where $E_r'(H)$ is the first derivative of the r'th energy level with respect to the field and $P_r(T)$, the probability of finding a number of the ensemble in state $|r\rangle$

$$\begin{aligned} \chi_T^{xx} &= \left(\frac{\partial M^x}{\partial H^x} \right)_T = - \sum_r P_r E_r''(H^x) - \sum_r \left(\frac{\partial P_r}{\partial H} \right)_T E_r'(H^x) \\ &= \chi(\omega=0) - \sum_r \left(\frac{\partial P_r}{\partial H} \right)_T E_r'(H^x) \end{aligned} \quad (5.1.1)$$

and $E_r'(H^x = 0)$ is not zero when there is degeneracy, that is different spin configurations with same energy.

Allan and Betts (1968) pointed out that spins which have an equal number of neighbours pointing up as down, are essentially free. This causes the degeneracy mentioned above. We want to note here that for the XY model in any lattice even 1 dimensional there is no such degeneracy when we deal with the magnetization in the x direction, which is parallel to the interaction. As discussed in the next section the magnetization in the z direction is non-ergodic at least in the 1-dimensional case.

For the XY model we had to use series expansion methods for the time, or frequency, variable as well as for the temperature while the Ising model results mentioned above had closed formulae for the time variable. Series expansions were only needed for the static spin correlations of the spin at the origin with k of its nearest neighbours, $k \leq q$.

Another subject we want to mention is the critical slowing down for dynamical susceptibility. Suzuki (1969 a,b) calculated the dynamical susceptibility for the kinetic Ising model in two dimensions (and for the XY chain which we discuss in the next section). We quote

here some of his results as this was the only calculation of a dynamic susceptibility for a physically plausible model before our work, and we use his definitions and terminology.

The kinetic Ising model has transition probabilities defined by

$$\omega_j(\sigma_o) = \frac{1}{2\tau} (1 - \sigma_j \tanh \beta \sum_k J_{jk} \sigma_k) . \quad (5.1.2)$$

τ indicates the relaxation time of a free spin interacting with a heat bath. This system has a Louville operator

$$\mathcal{L} = \sum_k \omega_k(\sigma_k) (1 - P_k)$$

where P_k flips the k 'th spin. The dynamical susceptibility can be expanded in a moment series:

$$\chi(\omega) = N \mu m^2 \left\{ \frac{\mu_1}{i\omega\tau} + \frac{\mu_2}{(i\omega\tau)^2} + \dots \right\} \quad (5.1.3)$$

where

$$\mu_n = \frac{\tau^n}{Nm^2} \langle M \mathcal{L}^n M \rangle . \quad (5.1.4)$$

All moments μ_n are positive and remain finite at the critical point. Suzuki proved

$$\frac{Nm^2\beta\mu_2}{(\omega\tau)^2 + \frac{\mu_4}{\mu_2}} \leq \operatorname{Re} \chi(\omega) \leq \frac{Nm^2\beta\mu_2}{(\omega\tau)^2} \quad (5.1.5)$$

and therefore postulates that the susceptibility behaves like

$$\chi(\omega) \sim \frac{1}{i\omega + \epsilon^\gamma f\left(\frac{i\omega}{\epsilon^\gamma}, \epsilon\right)} \quad (5.1.6)$$

and $f(0, \epsilon)$ can be singular at $\epsilon=0$. Expanding in powers of ω

$$\chi(\omega) = N\beta \sum (-i\omega\tau)^n a_n(\epsilon) \quad (5.1.7)$$

and calculation gave

$$a_n(\epsilon) \sim \epsilon^{-\alpha_n} \quad (5.1.8)$$

$$\alpha_n \geq \gamma + n \Delta_s$$

where

$$\Delta_s = \alpha_1 - \gamma \quad (5.1.9)$$

and $\Delta_s \geq \gamma$ was proved by Abe and Hatano (1969). The index of critical slowing down Δ_s is defined from

$$\tau_M = \int_0^\infty \frac{\langle M(t)M \rangle}{\langle M^2 \rangle} dt \sim \epsilon^{-\Delta_s} \quad (5.1.10)$$

and is found to be the same as in (5.1.9). Suzuki defines the index of critical slowing down of the self correlation (see 5.4)

$$\int_0^{\infty} \langle \sigma_1(t) \sigma_1 \rangle dt \sim \epsilon^{-\Delta_A} . \quad (5.1.11)$$

Classical theories predict $\Delta_s = \gamma$. Suzuki defines $R_M(\epsilon)$ by

$$\tau_M = \chi(0) R_M^{-1}(\epsilon) \quad (5.1.12)$$

and states a proposed law of similarity as follows. The relaxation time singularity consists of the direct or thermodynamic critical slowing down due to the susceptibility and the indirect, due to $R(\epsilon)$, called induced critical slow down (or speed up as the case may be). Relaxation times can be defined for various quantities not only the magnetization. The similarity law states that the indirect critical slow down is the same for some critical variables. Variables for which it holds are called similar variables.

For the kinetic Ising model in 2 dimensions Suzuki expects that the magnetization and the energy are similar variables.

$$\tau_E = \langle \delta E \frac{1}{\mathcal{L}} \delta E \rangle / \langle (\delta E)^2 \rangle \sim \epsilon^{-\Delta_E} \quad (5.1.13)$$

$$\tau_{EM} = \langle \delta E \frac{1}{\mathcal{L}} M \rangle / \langle M \delta E \rangle \sim |\epsilon|^{-\Delta_{ME}} \quad (5.1.14)$$

$$\tau_M = \langle M \frac{1}{\mathcal{L}} M \rangle / \langle M^2 \rangle \sim \epsilon^{-\Delta_M} \quad (5.1.15)$$

and we know $\langle (\delta E)^2 \rangle \sim \epsilon^{-\alpha}$ and $\langle M \delta E \rangle \sim |\epsilon|^{\beta-1}$. Only Δ_M is calculated directly but Suzuki using the relation

$$\langle (\lambda M + \delta E) \frac{1}{\mathcal{L}} (\lambda M + \delta E) \rangle \geq 0 \quad (5.1.16)$$

proves that

$$2\Delta_{ME} - (\Delta_M + \Delta_E) \leq \alpha + 2\beta + \gamma - 2 \quad (5.1.17)$$

for real λ . Similarity and (5.1.17) yield $\alpha + 2\beta + \gamma \geq 2$ which is valid assuming scaling. Suzuki obtains that

$\Delta_M = 2.00 \pm 0.05$. Similarity claims $\Delta_E - \alpha = \Delta_M - \gamma$ and since $\alpha = 0$ $\gamma = \frac{7}{4}$ this yields $\Delta_E = \frac{1}{4}$. This is consistent with computer simulations.

We calculated a series for $\chi^{xx}(\omega)$ by methods explained in sections (5.4) and (5.5) and estimated the critical index of slowing down Δ_s obtaining that in this case it is actually a speeding up, $\Delta_s < \gamma$.

and $g^+(\infty, 0) = g^-(\infty, 0)$. $y = d+x$ where d is the dimensionality. The indices η and ν were defined in (2.1.8) and (2.1.9) respectively. The Fourier transform of (2.1.9) is

$$\Gamma(\underline{k}, \varepsilon=0) \propto k^{y-z} \quad (5.2.5)$$

$$\Gamma(\underline{k}=0, \varepsilon) \propto \varepsilon^{-Y} \propto \xi^{Y/\nu} \quad (5.2.6)$$

as $\xi \sim \varepsilon^{-1/\nu}$. From (5.2.4) we see

$$\Gamma(\underline{k}, \xi) \propto \xi^{-Y} [1 + \dots] \quad \text{for } k\xi \ll 1 \quad (5.2.7)$$

if $\Gamma(\underline{k}=0, \varepsilon)$ is finite for $\frac{1}{\xi} \neq 0$. Similarly if $\Gamma(k, \varepsilon=0)$ is finite for $k \neq 0$

$$\Gamma(\underline{k}, \xi) \propto k^Y [1 + \dots] \quad k\xi \gg 1. \quad (5.2.8)$$

Comparing (5.2.5), (5.2.6) with (5.2.7) and (5.2.8), the scaling law

$$y = -2 + \eta = -Y/\nu \quad (5.2.9)$$

results. Adding assumptions on the structure of $\Gamma(\underline{k}, \xi, H)$, all scaling laws (2.4.24) predicted by the Kadanoff construction are obtained. Returning to $\Gamma(\underline{r}, \xi)$, since

$$\mathbf{x} = \mathbf{y} - \mathbf{d}$$

$$\Gamma(\underline{r}, \xi) = r^{\eta-2-d} g(r/\xi) \quad (5.2.10)$$

The dynamical scaling assumption is a generalization of (5.2.4) to the frequency dependent correlation. $\Gamma_{\xi}(\underline{k}, \omega)$ is defined by

$$\Gamma_{\xi}(\underline{r}, t) = \int \frac{d^3 \underline{k}}{(2\pi)^3} \int \frac{d\omega}{2\pi} e^{i(\underline{k} \cdot \underline{r} - \omega t)} \Gamma_{\xi}(\underline{k}, \omega) \quad (5.2.11)$$

where

$$\Gamma_{\xi}(\underline{r}, t) = \langle \{S(\underline{r}, t) - \langle S(\underline{r}, t) \rangle, S(0,0) - \langle S(0,0) \rangle\} \rangle \quad (5.2.12)$$

The curly brackets are anticommutator. ξ , the correlation length, gives the dependence on the temperature.

The scaling assumption is

$$\Gamma_{\xi}(\underline{k}, \omega) = \frac{2\pi \Gamma_{\xi}(\underline{k}) f_{k\xi}\left(\frac{\omega}{k^z \Omega(k\xi)}\right)}{k^z \Omega(k\xi)} \quad (5.2.13)$$

where

$$\Gamma_{\xi}(\underline{k}) = \int_{-\infty}^{\infty} \frac{d\omega}{2\pi} \Gamma_{\xi}(\underline{k}, \omega) \quad (5.2.14)$$

and (5.2.14) implies then

$$\int_{-\infty}^{\infty} f_{k\xi}(x) dx = 1 \quad (5.2.15)$$

Ω is determined from the constraint

$$\int_{-1}^1 f_{k\xi}(x) dx = \frac{1}{2} . \quad (5.2.16)$$

The function $k^z \Omega(k\xi)$ is called the characteristic frequency $\omega_{\xi}^M(\underline{k})$. The time dependent correlation is scaled as

$$\Gamma_{\xi}(\underline{k}, t) \propto \frac{\Gamma_{\xi}(\underline{k}) \hat{f}(t.k^z \Omega(k\xi))}{k^z \Omega(k\xi)} . \quad (5.2.17)$$

For the isotropic Heisenberg model the Landau-Lifshitz (1935) formula gives

$$\omega_{\xi}^M(\underline{k}) = \lambda k^2$$

where

$$\lambda = \frac{\rho_s}{| \langle M \rangle |} ,$$

and ρ_s is a stiffness constant proportional to ξ^{-1} .

From this Halperin and Hohenberg deduce to first order

$$z = 3 - \beta/\nu .$$

z cannot be determined from thermodynamic considerations in the anisotropic case.

We calculated $\langle \sigma_1(t) \sigma_1 \rangle$, the self correlation.
The NMR line width Δ_{NMR} is given by (Heller (1966)):

$$\Delta_{\text{NMR}} \propto \int_0^{\infty} \langle \sigma_1(t) \sigma_1 \rangle dt \quad (5.2.18)$$

$$\propto \int_0^{\infty} \Gamma_{\xi}(r=0, t) dt$$

$$\propto \int_0^{\infty} d^3k \Gamma_{\xi}(\underline{k}, \omega=0)$$

$$\propto \int_0^{\infty} \frac{k^2 dk \Gamma_{\xi}(\underline{k}) f_{k\xi}(0)}{k^z \Omega(k\xi)}$$

$$\propto \int_0^{\infty} dk k^{2-z+\eta-2} \frac{g(k\xi) f_{k\xi}(0)}{\Omega(k\xi)}$$

$$\propto \xi^{z-\eta-1} \int_0^{\infty} dx F(x) \quad .$$

Since $\xi \sim \epsilon^{-\nu}$

$$\Delta_{\text{NMR}} \propto \epsilon^{-\nu(z-\eta-1)} \quad . \quad (5.2.19)$$

The assumption that (5.1.13) holds for the order parameter is restricted scaling. Extended scaling is when (5.2.13) holds for other operators also.

Assuming that the frequency scales like $\omega/\epsilon^{\Delta_S}$ yields, using (5.2.13), that

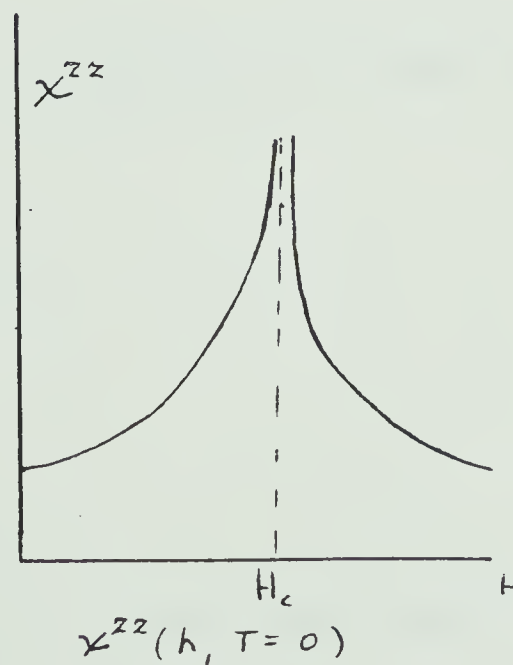
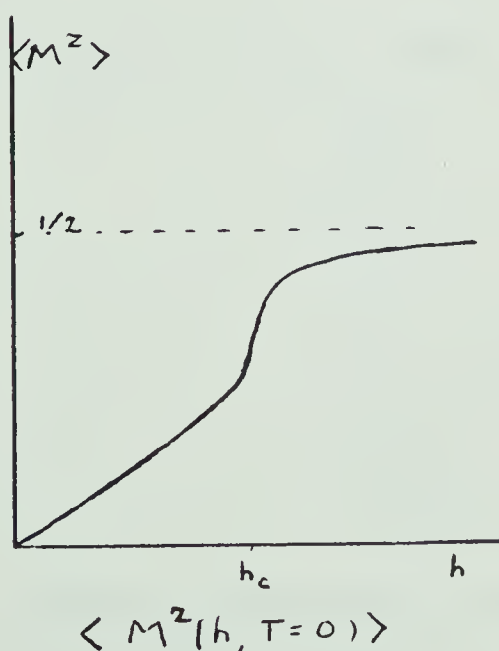
$$z = \nu^{-1} \Delta_S \quad . \quad (5.2.20)$$

5.3 Dynamical results for the one dimensional XY model

Many authors have studied the time dependent correlations in the one dimensional XY model. (a) Niemejer (1967, 1968), (b) Katsura, Horiguchi and Suzuki (1969), (c) Suzuki (1969), and (d) Barouch, McCoy and Dresden (1970).

(a) and (d) were concerned with the exact solution of the anisotropic XY model which can be obtained for arbitrary fields in the z direction. Of course specific cases and asymptotic limits had to be taken to see what those solutions mean. (b) obtained linear response for the isotropic XY model by the two time Green's function method. (c) discussed the results for the anisotropic XY model terms of critical slowing down.

In (a) the static result obtained was that the one dimensional XY model has a divergence at $T=0$, $h=h_c$,



This makes the order in which limits are taken to obtain asymptotic behaviour crucial . The susceptibility behaves like

$$\chi_{T=0}(h) \approx \log (|h-1|) \quad , \quad (5.3.1)$$

where

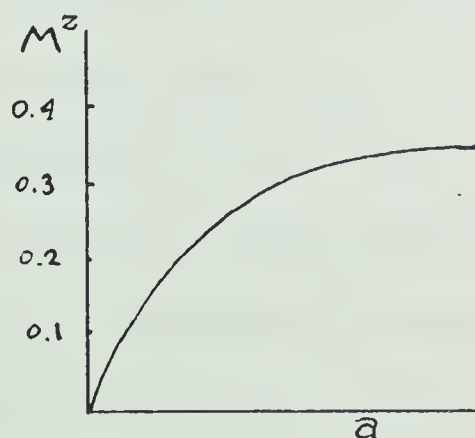
$$h = \frac{u(H-H_c)}{J\gamma}$$

γ being the anisotropy parameter. The points of interest to us were in (d) and (c).

In (d) the authors found that the asymptotic behaviour of $M(t)$ is the same for large times for all H , only the coefficients depending on the details. There is no approach to equilibrium no matter how slowly the field varies. The easiest example they provide is

$$H(t) = \begin{cases} a & t \leq 0 \\ 0 & t > 0 \end{cases} .$$

The plot of $M_z(t=\infty, T=0)$ as a function of a looks like



with M_z definitely saturating at a value less than $\frac{1}{2}$,

where $\frac{1}{2}$ is the saturation value of the initial magnetization.

Suzuki obtains that the dynamical susceptibility behaves like

$$\text{Re } \chi(\omega, h_c, T=0) \sim -\log|\omega| \quad (5.3.2)$$

and despite lack of ergodicity

$$\lim_{t \rightarrow \infty} (M, M(t)) = \langle M \rangle^2$$

so he can define the relaxation time and obtains

$$\tau_M^2 = -\int_0^\infty \frac{(\delta M, \delta M(t))t}{(M, M)} dt \sim h^{-2\Delta_M} \quad (5.3.3)$$

where $\Delta_M = 1$.

For the partial energy defined by

$$E = 4 \sum_i S_i^x S_{i+1}^x \quad (5.3.4)$$

he obtains $\Delta_E = 1$ and since $\gamma_E = 0$ similarity holds for this partial energy and the susceptibility.

We note that here $\Delta_M > \gamma$, but here because of the singular character of the model the temperature was kept fixed at $T=0$ and the magnetic field is the variable, which is quite different from our procedure. Also these results were obtained by Suzuki for the anisotropic XY model as $\text{Re } \chi(k=0, \omega) = 0$, according to his calculations, when the anisotropy vanishes.

5.4 The Kubo linear response theory

The Kubo formalism (Kubo 1957) is a general approach to the problem of time dependent disturbances. The results are perturbation series in the disturbance. We shall, as most users of the formalism do, keep only the term linear in the disturbance.

The system is assumed to be in equilibrium without disturbances at $t = -\infty$. Its properties are then described by a density matrix ρ_0 . For a canonical ensemble

$$\rho_0 = \frac{e^{-\beta \mathcal{H}_0}}{\text{Tr } e^{-\beta \mathcal{H}_0}} \quad (5.4.1)$$

A disturbance is introduced; the external force $F(t)$ is conjugate to the observable A . The perturbed Hamiltonian is now

$$\mathcal{H} = \mathcal{H}_0 + \mathcal{H}' = \mathcal{H}_0 - \vec{A} \cdot \vec{F}(t) \quad (5.4.2)$$

The time evolution of the system is given by the equation of motion of the density matrix

$$\frac{d}{dt} \rho(t) = \frac{1}{i\hbar} [\mathcal{H}, \rho(t)] \quad (5.4.3)$$

The solution of (5.4.3) after linearization is:

$$\rho(t) = \rho_0 + \frac{1}{i\hbar} \int_{-\infty}^t e^{-\frac{i(t-t')\mathcal{H}_0}{\hbar}} [\mathcal{H}, \rho_0] e^{\frac{i(t-t')\mathcal{H}_0}{\hbar}} dt \quad (5.4.4)$$

The response of an observable B to the disturbance is

$$\Delta B = \text{Tr} (\rho(t) - \rho_0) B .$$

ΔB is the macroscopic difference of the average of B at time t from its equilibrium average. It follows that:

$$\Delta B = \frac{1}{i\hbar} \text{Tr} \int_{-\infty}^t [\rho_0, A] B(t-t') F(t') dt' \quad (5.4.5)$$

where

$$B(t) = e^{\frac{it}{\hbar} \mathcal{H}_0} B e^{-\frac{it}{\hbar} \mathcal{H}_0} .$$

The linear response function ϕ_{BA} , which is the response of the observable B to a disturbance conjugate to A is defined by

$$\Delta B(t) = \int_{-\infty}^t \phi_{BA}(t-t') F(t') dt' . \quad (5.4.6)$$

It follows then, that

$$\phi_{BA}(t) = \frac{1}{i\hbar} \text{Tr} \rho_0 [A, B(t)] . \quad (5.4.7)$$

Many disturbances of interest can be expanded in Fourier series. We therefore seek the response to a harmonic force. It is convenient to take F as

$$F(t) = \text{Re } F_0 e^{i\omega t + \epsilon t} \quad (5.4.8)$$

where ε is a small positive factor which ensures that $F(-\infty)=0$ as an initial unperturbed condition. This is a slow quasistatic switching on of the force.

The frequency dependent susceptibility is defined by

$$\Delta B(t) = \text{Re } \chi_{BA}(\omega) F_0 e^{i\omega t} . \quad (5.4.9)$$

The result of combining (5.4.9) and (5.4.6) is:

$$\chi_{BA}(\omega) = \lim_{\delta \rightarrow 0+} \int_0^{\infty} \phi_{BA}(t) e^{-i\omega t - \delta t} dt . \quad (5.4.10)$$

The response function is real and the susceptibility satisfies the following symmetry relations

$$(1) \quad \text{Re } \chi_{BA}(\omega) = \text{Re } \chi_{BA}(-\omega)$$

$$(2) \quad \text{Im } \chi_{BA}(\omega) = -\text{Im } \chi_{BA}(-\omega)$$

$$(3) \quad \chi_{BA}(\omega, -F) = \varepsilon_A \varepsilon_B \chi_{AB}(\omega, F) ,$$

where $\varepsilon_A, \varepsilon_B$ are +1 or -1 as A and B are respectively even or odd with respect to time reversal. Relation (3) is known as the Onsager relation (1931).

For short times the response function can be expanded in an asymptotic power series. Using (3.3.17)

$$\begin{aligned} \phi_{BA}(t) &= \frac{1}{i\hbar} \text{Tr} \rho_0 \left[A, \left\{ B + \frac{it}{\hbar} [\mathcal{H}, B] + \left(\frac{it}{\hbar}\right)^2 [\mathcal{H}, [\mathcal{H}, B]] + \dots \right\} \right] \\ &= \sum f_n t^n \end{aligned} \quad (5.4.11)$$

This series is either odd or even in t and in our magnetic case ϕ_{MM} is odd. Substituting (5.4.11) in (5.4.9) and using

$$\int_0^{\infty} x^m e^{-\epsilon x} e^{-i\omega x} dx = \pi i^m \delta^m(\omega - i\epsilon) + \frac{m!}{(i\omega + \epsilon)^{m+1}}. \quad (5.4.12)$$

we obtain, for ϕ_{BA} odd in t ,

$$\text{Re } \chi_{BA}(\omega) = \sum \frac{m! (-1)^{\frac{m+1}{2}}}{\omega^{m+1}} f_m \quad (5.4.13)$$

$$\text{Im } \chi_{BA}(\omega) = \pi (i)^m \delta^m(\omega) f_m.$$

The initial values of the response function (since $f_m = \frac{1}{m!} \phi^{(m)}(0)$) determine the coefficients of the power series in $\frac{1}{\omega}$ for the susceptibility and the moments for the dissipative part. The quantity we are calculating here is $\chi^{xx}(\omega)$, that is the force we apply is $H^x e^{i\omega t}$ and both operators A and B are M^x . The response function is:

$$\phi_{xx}(t) = \frac{i}{\hbar} \text{Tr } \rho_0 [M_x, M_x(t)] \quad (5.4.14)$$

and the susceptibility follows from (5.4.11) and (5.4.13). Obviously in this case $\phi_{xx}(t) = -\phi_{xx}(-t)$. It can be noted, that

$$\phi_{xx}(t) = \phi_{yy}(t)$$

and

$$\begin{aligned}\phi_{xy}(t) &= \phi_{yx}(t) = \phi_{xz}(t) = \phi_{yz}(t) \\ &= \phi_{zx}(t) = \phi_{zy}(t) = 0.\end{aligned}$$

But for neutron scattering where one needs the $\chi(\underline{q}, \omega)$ the zz part does not vanish.

We want to obtain the zero frequency susceptibility. The identity

$$[A, e^{-\beta \mathcal{H}_0}] = \frac{\hbar}{i} e^{-\beta \mathcal{H}_0} \int_0^\beta \dot{A}(-i\hbar\lambda) d\lambda \quad (5.4.15)$$

was proved by Kubo (1957) and independently by Feynman (1948). Using (5.4.15) and integrating by parts we obtain the following expression in the zero frequency limit

$$\chi_{BA}(0) = \int_0^\beta \text{Tr } \rho_0 A(-i\hbar\lambda) B d\lambda - \beta \text{Tr } \rho_0 A^0 B^0. \quad (5.4.16)$$

$\chi_{BA}(0)$ is the quasistatic (also called adiabatic) susceptibility and

$$A^0 = \lim_{T \rightarrow \infty} \frac{1}{T} \int_0^T A(t) dt.$$

The isothermal susceptibility on the other hand is

$$\chi_{BA}^T(0) = \frac{\text{Tr } \rho B - \text{Tr } \rho_0 B}{F} \quad (5.4.17)$$

which can be shown (Kubo 1957) to be equal to

$$\chi_{BA}^T(0) = \int_0^\beta \text{Tr } \rho_O A(-i\hbar\lambda) B d\lambda - B\langle A \rangle \langle B \rangle \quad (5.4.18)$$

in the linear approximation.

The second term in (5.4.16) was obtained from

$$\lim_{t \rightarrow \infty} \int_0^\beta \text{Tr } \rho_O A(-i\hbar\lambda) B(t) d\lambda$$

by noting that if this limit exists it is equal to

$$\begin{aligned} \lim_{t \rightarrow \infty} \frac{1}{t} \int_0^t dt \int_0^\beta \text{Tr } \rho_O A(-i\hbar\lambda) B(t) d\lambda &= \\ &= \int_0^\beta \text{Tr } \rho_O A(-i\hbar\lambda) B_O d\lambda . \end{aligned}$$

This limiting process picks up the zero frequency component of B , B_O , which is diagonal with respect to \mathcal{H}_O .

Now

$$\begin{aligned} \text{Tr } \rho_O Q B_O &= \sum_{i,j,k} \rho_O^{ij} Q^{jk} B_O^{ki} = \\ &= \sum_i \rho_O^{ii} Q^{ii} B_O^{ii} = \text{Tr } \rho_O Q_O B_O \end{aligned}$$

and the diagonal part of $A(-i\hbar\lambda)$ is just A^O since $(e^{-\lambda\mathcal{H}} A e^{\lambda\mathcal{H}})_{ii} = A_{ii}$.

For the isothermal and adiabatic susceptibility to be equal it is sufficient that

$$\lim_{t \rightarrow \infty} \langle AB(t) \rangle = \langle A \rangle \langle B \rangle \quad (5.4.19)$$

A system which satisfies (5.4.19) is called ergodic. For a discussion of the ergodicity of our system, see Section (5.1).

5.5 The high temperature series expansion for the frequency dependent susceptibility

As we saw in the previous section, the susceptibility can be expanded in an asymptotic series in powers of ω^{-1} . The coefficients of this expansion are determined by the time derivatives of the response function at zero time.

Since

$$\langle M_x(t) \rangle = \int_0^t \phi_{xx}(t-t') H_x(t') dt' \quad (5.5.1)$$

we expect ϕ_{xx} to be linear in N , the number of lattice sites. Therefore in the terminology of Sec. 3.2

$$\begin{aligned} \phi_{xx} &= \{ \text{linear part in } N \text{ of } -\frac{i}{\hbar} \langle [M_x, M_x(t)] \rangle \} \\ &= -\frac{i}{\hbar} \text{Tr} e^{KP} [M_x, M_x(t)] \quad (5.5.2) \end{aligned}$$

Defining $\omega_0 = J/\hbar$ we have

$$\phi_{xx}(t) = -\frac{i}{\hbar} \text{Tr}^* e^{KP} [M_x, e^{-iP\omega_0 t} M_x e^{iP\omega_0 t}] \quad (5.5.3)$$

We expand all exponentials in (5.5.3) in power series of P and have a double series in K and t

$$\begin{aligned} \phi_{xx} &= -\frac{im^2}{4\hbar} \text{Tr}^* \sum_{m,n,q=0}^{\infty} \frac{K^n}{n!} \frac{(i\omega_0 t)^{m+q}}{m! q!} \times \\ &\times \sum_{i,j} (P_{\sigma_i}^n (-P)_{\sigma_j}^m P^q - P^n (-P)_{\sigma_i}^m P^q_{\sigma_j}) \quad (5.5.4) \end{aligned}$$

where m' is the magnetic moment of the atoms.

We rearrange the sum, use the cyclic property of the trace and interchange m and q in the first term of (5.5.4) to obtain:

$$\begin{aligned} \phi_{xx} = & -\frac{im'^2}{4\hbar} \text{Tr}^* \sum_{n,s} \frac{K^n}{n!} \frac{(i\omega_0 t)^s}{s!} \times \\ & \times \sum_{m=0}^s \binom{s}{m} \sum_{i,j} P^{n+m} \sigma_i^x P^{s-m} \sigma_j^x ((-1)^{s-m} - (-1)^m). \end{aligned} \quad (5.5.5)$$

For even s $[(-1)^{s-m} - (-1)^m]$ vanishes. Therefore,

$$\begin{aligned} \phi_{xx} = & \frac{im'^2}{2\hbar} \text{Tr}^* \sum_n \frac{K^n}{n!} \sum_s \text{odd} \frac{(i\omega_0 t)^s}{s!} \times \\ & \times \sum_{m=0}^s \binom{s}{m} (-1)^m \sum_{i,j} P^{n+m} \sigma_i^x P^{s-m} \sigma_j^x. \end{aligned} \quad (5.5.6)$$

Substituting $\sigma_i^x = a_i^\dagger + a_i$

$$\begin{aligned} \phi_{xx}(t) = & \frac{im'^2}{2\hbar} \sum_{n=0}^{\infty} \frac{K^n}{n!} \sum_s \text{odd} \frac{(i\omega_0 t)^s}{s!} \times \\ & \times \sum_{m=0}^s \binom{s}{m} (-1)^m \text{Tr}^* \left[\sum_{i,j} (P^{n+m} a_i^\dagger P^{s-m} a_j + P^{n+m} a_i P^{s-m} a_j^\dagger) \right]. \end{aligned} \quad (5.5.7)$$

We have in (5.5.7) an expression similar to the one we had for the static susceptibility (3.3.10) with an added complication. Now a different factor multiplies different

vertical orderings of the arrows.

Instead of one vertical weight we now calculate $\ell+1$ weights, where ℓ is the number of solid arrows in the graph. As before the tip of the dotted $i \rightarrow j$ arrow is fixed - at the zero level. The tail is fixed in sequence at each of the levels k and the number of allowed orderings of the other arrows is $v(k-2)$. For a fixed n and s

$$\begin{aligned} \sum_{m=0}^s \binom{s}{m} (-1)^m \text{Tr}^* P^{n+m} a_i^\dagger P^{s-m} a_j &= \\ = \sum_{g_{s+n+1}} \sum_{m=0}^s \binom{s}{m} (-1)^m \frac{(g_{s+n}, L) h(g') v(s-m)}{2^v} &\quad (5.5.8) \end{aligned}$$

where v is the number of vertices.

As in the susceptibility the graphs were either "bubbles" for $i=j$, or fluctuation type graphs with a dotted arrow $i-j$ unrestricted in length. It can be noted though that when $v(k)=v$ for all k , we have in (5.5.9)

$$\sum_{m=0}^s \binom{s}{m} (-1)^m v = 0$$

so graphs where the $i-j$ arrow has a vertex of order 2 do not contribute, which reduces greatly the number of

fluctuation type graphs. There is no need to calculate all $v(k)$, since using the cyclic property of the trace we can obtain that

$$v(k) = v(\ell - k) \quad . \quad (5.5.9)$$

The equivalent of Theorem I is a bit more complicated here.

Theorem Ia

Let g be a disconnected graph of ℓ solid arrows and one dotted arrow. There are two subgraphs, g_1 and g_2 . g_1 has ℓ_1 solid arrows and a dotted one and g_2 has ℓ_2 solid arrows. Then:

$$v(k) = v_2 \sum_{i=0}^{k-1} v_1(i) \binom{k-1}{i} \binom{\ell-k+1}{\ell_1-i} \quad (5.5.10)$$

where $v_1(i)$ is the vertical weight of g_1 with the dotted tail at the i th level. $v(k)$ is the vertical weight for the whole graph g with tail of the dotted arrow at the k th level. v_2 is the vertical weight of g_2 as a partition function graph.

Proof

We keep the dotted tail fixed at the k th level. Among the $k-1$ levels between the tip and tail of the dotted arrows i will be filled by arrows from g_1 . The

number i can run from zero to $k-1$ and

$$v(k) = v_2 \sum_{i=0}^{k-1} v_1(i+1) X_i$$

where X_i is the number of ways ℓ arrows, $\ell_1 \in g_1$ and $\ell_2 \in g_2$, can be assigned to ℓ levels in such a way as to have i arrows belonging to g_1 assigned to the first $k-1$ levels.

The result (5.5.10) follows when we note

$$X_i = \frac{(k-1)!}{i!(k-1-i)!} \frac{(\ell-k+1)!}{(\ell_1-i)! (\ell_2-k+1-i)!} \cdot$$

The first factor is choosing i levels out of $k-1$ to fill with g_1 . The second factor is the number of ways of assigning $\ell-k+1$ levels so that ℓ_1-i are occupied by arrows of g_1 and the rest by arrows of g_2 ,

The equivalent of Theorem II holds too.

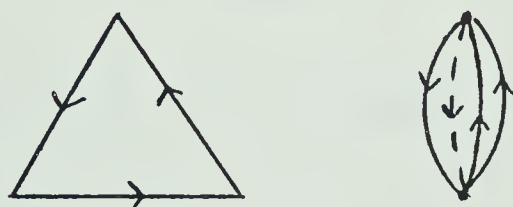
Theorem IIa

Let a connected graph g of ℓ solid arrows and a dotted arrow be composed of two subgraphs, g_1 of ℓ_1 solid arrows and the dotted arrow and g_2 of ℓ_2 solid arrows such that g_1 and g_2 have m common vertices of order 2 in g and no other common vertices. Then the vertical weights $v(k)$ of g are given by the same expression (5.5.10) as in Theorem Ia.

Proof

Exactly as the proof of Theorem II followed that of Theorem I. We have no restrictions on vertices of order 2, so they can be cut to give two separated graphs.

A numerical example to clarify the meaning of the theorem: The graph g is composed of g_1 and g_2



$\ell = 6$, $\ell_1 = \ell_2 = 3$. v_2 , the vertical weight of g_2 alone, is $3! = 6$. $v_1(i)$ are the vertical weights for g_1 where $i-1$ levels separate head and tail of the dotted arrow. Obviously $v_1(1) = v_1(4) = 1$, $v_1(2) = v_1(3) = 0$. Let us calculate $v(3)$, the vertical weight of g when 2 levels separate head and tail of dotted arrow.

Substituting in (5.5.10)

$$v(3) = 6 \cdot 1 \cdot \binom{2}{0} \binom{4}{3} = 24 .$$

With some labour the same result can be obtained directly.

Theorem IIIa

Let a connected graph g of ℓ solid arrows and a dotted arrow be composed of two subgraphs g_1 and g_2 of

ℓ_1 and ℓ_2 solid arrows respectively and a third subgraph consisting of the single dotted arrow. The subgraphs g_1 and g_2 have in common m vertices of order 2 in g and no other common vertices. Then the vertical weights are given by:

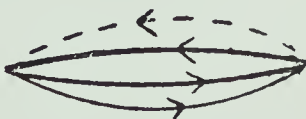
$$v(k) = \sum_{i=0}^k v_1(i+1) v_2(k-i) \binom{k-1}{i} \binom{\ell-k+1}{\ell_1-i} \quad (5.5.11)$$

The dotted arrow belongs to both g_1 and g_2 for the purpose of calculating the $v_1(i+1)$, $v_2(k-i)$. The proof follows exactly the proof of Ia. After assigning i levels from the first $k-1$ levels to g_1 , we assign the $(k-1-i)$ levels left to g_2 . The dotted arrow belongs to both, so the vertical weights of both are dependent on i .

In the actual calculation we did not go far enough to need the theorems, their only use was for checking the computer programs.

Using the fact that vertices of order 2 introduce no added restriction the dotted arrow was broken by adding a vertex of order 2 and the two new arrows were put at the head of the vector describing the graph.

In Appendix I the graph



is described by the vector $13 \ 32 \ 21 \ 12 \ 21 \ .$

Thus the first arrow could be kept fixed at the zeroth level and the second at the k 'th with no difficulty.

5.6 The frequency dependent susceptibility series, analysis and conclusions

The series we obtained for the frequency dependent susceptibility on the f.c.c. lattice is:

$$\chi(\omega) = \frac{-m'}{2\hbar} \sum_{s \text{ odd}} \left[s! \left(\frac{\omega_0}{\omega} \right)^{s+1} \chi_{s+1}(K) \right]$$

$$\chi_2(K) = 3K + 9K^2 + 14.75K^3 + 22.0K^4 + 61.375K^5 + 237.8041667K^6 + \\ .880.980209K^7 + O(K^8)$$

$$\chi_4(K) = 20K + 6K^2 - 149.5K^3 - 575.291667K^4 - 1433.55K^5 + O(K^6)$$

$$\chi_6(K) = 91.5K + 83.825K^2 - 296.1208K^3 + O(K^4)$$

$$\chi_8(K) = 391.516667K + O(K^2)$$

As part of the calculation we had to calculate $\langle [\sigma_1(t), \sigma_1] \rangle$, the bubble graphs. As can be seen from (5.5.5) the odd powers of t for $\langle [\sigma_1(t), \sigma_1] \rangle$ are just twice those of $\langle \sigma_1(t) \sigma_1 \rangle$, so we have the coefficients of the first, third, fifth and seventh powers of t .

We also calculate the coefficients of the second and fourth powers of t in order to see how the dependency on K changes with the powers of t , and estimate $\int_0^\infty \langle \sigma_1(t) \sigma_1 \rangle dt$ (see (5.2.18)). The results are

$$\langle \sigma_1(t) \sigma_1 \rangle = 1 + \sum_{i=1}^{\infty} a_i(K) (-i\omega t)^i$$

$$a_1 = 3K + 6K^2 + 10.75K^3 + 22.5K^4 + 64.175K^5 + 218.691667K^6 + \\ + 780.8421131K^7 + O(K^8)$$

$$a_2 = 6 + 12K + 49.5K^2 + 142.0K^3 + 374.25K^4 + 1054.65K^5 + 3317.85625K^6 + O(K^7)$$

$$a_3 = 28K + 97K^2 + 298K^3 + 853.166667K^4 + 2453.48333K^5 + O(K^6)$$

$$a_4 = 336 + 1164K + 6702K^2 + 24392K^3 + O(K^4)$$

$$a_5 = 156.3K + 707.7K^2 + 2756.3125K^3 + O(K^4)$$

$$a_7 = 664.37857K + O(K^2)$$

We assumed that $\langle \sigma_1(t) \sigma_1 \rangle$ is independent of N and therefore

$$\langle \sigma_1(t) \sigma_1 \rangle = \frac{\text{Tr } e^{-\beta \mathcal{H}} \sigma_1(t) \sigma_1}{\text{Tr } e^{-\beta \mathcal{H}}} = \frac{A(K, t) + NB(K, t) + \dots}{1 + a(K) N + \dots} \\ = A(K, t) \quad .$$

A. Analysis of second and fourth moments of the susceptibility

We did the regular Pade approximant analysis and ratio analysis as seen in the following tables.

TABLE 5.1

Pade approximants to $\frac{\partial}{\partial K} \log \chi_2(K)$

N	M = 1	M = 2	M = 3	M = 4
1	.1754	.1138	.2174	.2332
2	.1145	.1479	.2380	
3	.2292	.2547		
4	.2497			

TABLE 5.2

Pade approximants to $(K-K_c) \frac{\partial}{\partial K} \log \chi_2(K) \Big|_{K=.221}$

N	M = 0	M = 1	M = 2	M = 3	M = 4
1	.0407	.0511	.0911	.0925	.0870
2	.0513	.0377	.0926	.0914	
3	.0987	.0957	.0874		
4	.0954	.1006			
5	.0845				

TABLE 5.3

Ratios of $\chi_2(K)$

n	$r_n(\chi_2)$
1	
2	3
3	1.6389
4	1.4915
5	2.7898
6	3.8746
7	3.7046

TABLE 5.4

Pade approximants to $\frac{\partial}{\partial K} \log [(\chi_4(K))^{-1}]$

N	M = 1	M = 2
1	.1891	.2221
2	.2169	

TABLE 5.5

Pade approximants to $(K-K_c) \frac{\partial}{\partial K} \log \chi_4(K)^{-1} \Big|_{K=.221}$

N	M = 0	M = 1	M = 2
1	.735	.881	.877
2	.859	.877	
3	.875		

TABLE 5.6

Pade approximants to $(\chi_4(K)^{-1})^{1/0.8}$

N	M = 1	M = 2	M = 3	M = 4
0	*	.3481	.2443	.2235
1	*	.1814	.2111	
2	.3199	.2236		
3	.2014			

TABLE 5.7

Pade approximants $(\chi_4(K)^{-1})^{1/0.9}$

N	M = 1	M = 2	M = 3	M = 4
0		.3678	.2561	.2323
1	*	.1871	.2167	
2	.3154	.2282		
3	.2093			

B. Analysis of the coefficients of the first, second and third powers of t in $\langle \sigma_1(t) \sigma_1 \rangle$

TABLE 5.8

Pade approximants to $\frac{\partial}{\partial K} \log a_1(K)$

N	M = 1	M = 2	M = 3	M = 4	M = 5
0	.632	.414	.312	.276	.260
1	.352	.176	.232	.240	
2	.218	.225	.242		
3	.225	.211			
4	.239				
+					

TABLE 5.9

Pade approximants to $(K-K_c) \frac{\partial}{\partial K} \log a_1(K) \Big|_{K=0.221}$

N	M = 0	M = 1	M = 2	M = 3	M = 4	M = 5
0	.4420	.2679	.2007	.1672	.1471	.1333
1	.1547	.0827	.1007	.0955	.0852	
2	.0971	.0983	.0968	.1075		
3	.0983	.0976	.0988			
4	.0966	.0989				
5	.0893					

TABLE 5.10

Pade approximants to $\frac{\partial}{\partial K} \log a_2(K)$

N	M = 1	M = 2	M = 3	M = 4	M = 5
0			.2131	.2237 $\pm .1246i$.2401
1	.4237	*	.3333	.3061	
2	.7367	.5149	.2922		
3	.2970	.1277			
4	.1518				

TABLE 5.11

Pade approximants to $(K-K_c) \frac{\partial}{\partial K} \log a_2(K) \Big|_{K=0.221}$

N	M = 2	M = 3	M = 4	M = 5
0	.3102	.3226	.2475	.2322
1	.3218	.3116	.2268	
2	.1629	.0888		
3	.1154			

TABLE 5.12

Pade approximants to $(a_2(K))^{1/0.31}$

N	M = 1	M = 2	M = 3	M = 4	M = 5
1	.1575	.1548	.1894 $\pm .0735i$.1979 $\pm .0724i$.2054
2	.1984	.2736 $\pm .0495i$.2612 $\pm .0601i$.2803	
3	.2216	.2604 $\pm .0604i$.2695 $\pm .0545i$		
4	.2389	.2929 $\pm .0346i$			
5	.2507				

TABLE 5.13

Pade approximants to $(a_2(K))^{1/.25}$

N	M = 1	M = 2	M = 3	M = 4	M = 5
1	.140	*	.170 \pm .068i		.192
2	.176	.224 \pm .062i		.254	
3	.198	.224 \pm .062i	.224 \pm .062i		
4	.216	.250 \pm .052i			
5	.230				

TABLE 5.14

Pade approximants to $(a_2(K))^{1/.23}$

N	M = 1	M = 2	M = 3	M = 4	M = 5
1	.134	*	.163± .065i	*	.186
2	.167	.208± .062i	.212± .061i	.238	
3	.190	.212± .061i	.207± .063i		
4	.207	.237± .053i			
5	.222				

TABLE 5.15

Pade approximants to $\frac{\partial}{\partial K} \log a_3(K)$

N	M = 1	M = 2	M = 3
0	.373	.421	.340
1	.415	.385	
2	.326		

TABLE 5.16

Pade approximants to $(K-K_c) \frac{\partial}{\partial K} \log a_3(K) \Big|_{K=0.221}$

N	M = 0	M = 1	M = 2	M = 3
0	.766	.544	.414	.341
1	.454	.207	.139	
2	.242	.118		
3	.164			

TABLE 5.17

Pade approximants to $(a_3(K))^{1/0.35}$

N	M = 1	M = 2	M = 3	M = 4
0	.101	.138± .094i	.221	.221
1	.159	.197± .088i	.221	
2	.198	.244± .075i		
3	.225			

TABLE 5.18

Pade approximants to $(a_3(K))^{1/.38}$

N	M = 1	M = 2	M = 3	M = 4
0	.110	.155± .100i	.267	.232
1	.170	.213± .001i	.230	
2	.209	.262± .074i		
3	.235			

TABLE 5.19

Ratios of $a_1(K)$, $a_2(K)$ and $a_3(K)$

n	r_{a_1}	r_{a_2}	r_{a_3}
1		2	
2	2	4.125	3.4643
3	1.7917	2.8687	3.0722
4	2.0930	2.6356	2.8630
5	2.8522	2.8180	2.8757
6	3.4077	3.1459	
7	3.5705		

TABLE 5.20

Pade approximants to $[(K-K_c)^{0.1} a_1(K) K^{-1}]_{K=.221}$

N	M = 0	M = 1	M = 2	M = 3	M = 4	M = 5	M = 6
0		3.920	3.743	3.724	3.726	3.724	3.715
1	3.462	3.757	3.721	3.726	3.725	3.728	
2	3.683	3.730	3.727	3.725	3.727		
3	3.722	3.727	3.710	3.728			
4	3.727	3.726	3.729				
5	3.725	3.727					
6	3.720						

TABLE 5.21

Pade approximants to $[(K-K_c)^{0.24} a_2(K)]_{K=.221}$

N	M = 0	M = 1	M = 2	M = 3	M = 4	M = 5	M = 6
0		5.223	6.596	6.495	5.984	6.022	6.134
1	5.020	*	6.500	6.619	6.019	5.964	
2	5.879	6.275	6.006	6.132	6.102		
3	6.150	6.114	6.110	6.109			
4	6.109	6.109	6.109				
5	6.109	*					
6	6.109						

TABLE 5.22

Pade approximants to $[(K-K_c)^{0.35} a_3(K) K^{-1}]_{K=0.221}$

N	M = 0	M = 1	M = 2	M = 3	M = 4
0		28.25	26.67	24.83	23.77
1	23.37	26.78	34.69	22.03	
2	25.65	25.60	22.31		
3	25.60	25.65			
4	24.70				

The results for the self correlation indicate that

$$\begin{aligned} \langle \sigma_1(t) \sigma_1 \rangle &= 1 + \sum_{n=1}^{\infty} A_{2n} (K - K_C)^{-\gamma_{2n}^A} (-i\omega_0 t)^{2n} \\ &+ \sum_{n=0}^{\infty} A_{2n+1} K(K - K_C)^{-\gamma_{2n+1}^A} (i\omega_0 t)^{2n+1} \end{aligned} \quad (5.6.1)$$

with

$$\begin{aligned} \gamma_1^A &= 0.099 \pm 0.005 & A_1 &= 0.82 \pm 0.1 \\ \gamma_2^A &= 0.24 \pm 0.03 & A_2 &= 6.1 \pm 0.5 \\ \gamma_3^A &= 0.35 \pm 0.03 & A_3 &= 5.3 \pm 2, \end{aligned}$$

where the A_n were obtained by taking Pade approximants of the function $[(K-K_C)^{\gamma_n^A} a_n(K)]$ evaluated at $K=K_C$ for n even and of $[(K-K_C)^{\gamma_n^A} a_n(K)K^{-1}]$ evaluated at $K=K_C$ for n odd.

The differences are $\Delta_A = \gamma_n^A - \gamma_{n-1}^A \simeq 0.10 \pm 0.05$

$$\begin{aligned} \langle \sigma_1(t) \sigma_1 \rangle &\simeq 1 + \sum A_n (T-T_C)^{-n\Delta_A} (-i\omega_0 t)^n \\ &\simeq f\left(\frac{t}{\epsilon \Delta_A}\right) \end{aligned} \quad (5.6.2)$$

Near T_C , if we can assume that $\int_0^{\infty} f(x)dx$ converges, we find that the autocorrelation will vanish like ϵ^{Δ_A} .

$$\langle \sigma_1(t) \sigma_1 \rangle \simeq \int_0^{\infty} f\left(\frac{t}{\epsilon \Delta_A}\right) dt = \epsilon^{\Delta_A} \int f(x) dx \simeq \epsilon^{\Delta_A} \quad (5.6.3)$$

$$\text{We saw that } \Delta_{\text{NMR}} \propto \int_0^{\infty} \langle \sigma_1(t) \sigma_1 \rangle dt \simeq \epsilon^{\Delta_A} \quad (5.6.4)$$

and in (5.2.18) we obtained that $\Delta_{\text{NMR}} \simeq \epsilon^{v(\eta+1-z)}$,

in our case $\eta = 0$, $v = \frac{2}{3}$,

therefore $z = 0.85 \pm 0.10$. As seen in (5.2.20), $z = \frac{\Delta_s}{\nu}$; therefore $\Delta_s = 0.57 \pm 0.10$, where Δ_s is the critical slowing down index.

The result for Δ_A is corroborated by the calculation of Δ_s from the frequency dependent susceptibility and then substituting in (5.2.18) to obtain Δ_A .

As we know

$$\chi(\omega=0, \epsilon) \sim \epsilon^{-\gamma}.$$

The behaviour in ω we shall assume to be

$$\chi(\omega, \epsilon=0) \sim \omega^{-\gamma/\Delta_s}, \quad (5.6.5)$$

where this defines Δ_s .

A possible function with this behaviour and only even moments is

$$\chi(\omega) \sim \left[-\omega^2 + \epsilon^{2\Delta_s} f(\omega, \epsilon) \right]^{-\gamma/2\Delta_s} \quad (5.6.6)$$

and $f(\omega, \epsilon)$ regular at $\omega=0, \epsilon=0$.

The index of critical slowing down defined by Suzuki (1969) as mentioned in (5.1.10) is

$$\begin{aligned} \tau_2 &= \lim_{\omega \rightarrow 0} \left[\frac{\chi(\omega) - \chi(0)}{\omega^2 \chi(0)} \right]^{\frac{1}{2}} \\ &= \lim_{\omega \rightarrow 0} \left[\frac{(1-f)^{-1} \frac{\omega^2}{\epsilon^{2\Delta_s}} - 1}{\omega^2} \right]^{\frac{1}{2}} = \epsilon^{-\Delta_s} \end{aligned} \quad (5.6.7)$$

Our assumption is that $\chi(\omega)$ scales like $f\left(\frac{\omega}{\Delta_S}\right)$,
therefore the moments

$$\chi_2 \sim \epsilon^{-\gamma_2} = \epsilon^{-(\gamma-2\Delta_S)}$$

$$\chi_4 \sim \epsilon^{-\gamma_4} = \epsilon^{-(\gamma-4\Delta_S)}.$$

From our results

$$\gamma_2 = +.096 \pm .005 \rightarrow \Delta_S = .62 \pm .10$$

$$\gamma_4 = -.88 \pm .05 \rightarrow \Delta_S = .55 \pm .10.$$

This agrees with our former estimate.

Clearly our estimates of Δ_S however large the error margins yield

$$\Delta_S < \gamma$$

since

$$\Delta_S = 0.58 \pm 0.10.$$

This agrees with Tomita (1968) for the isotropic Heisenberg model where kinetic speeding up occurs, while the kinetic Ising model had $\Delta_S > \gamma$, a kinetic slowing down.

APPENDIX A

Graphs and Weights

In the following tables the graphs needed for the calculation of the static initial susceptibility, dynamical susceptibility and fourth order fluctuation are listed.

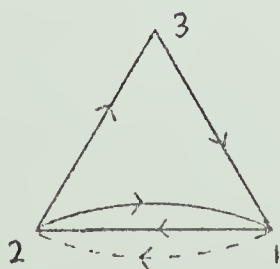
Table A.5 lists the fluctuation-like graphs used for the static susceptibility. The bubble graphs are not listed as they are identical to those of the dynamical susceptibility listed in Table A.7 apart from the vertical weight. The static weight is the sum of all dynamical weights. Table A.9 has the partition function graphs.

Table A.6 lists the fluctuation-like graphs generated for the dynamical susceptibility.

Table A.8 lists the fourth order fluctuation graphs.

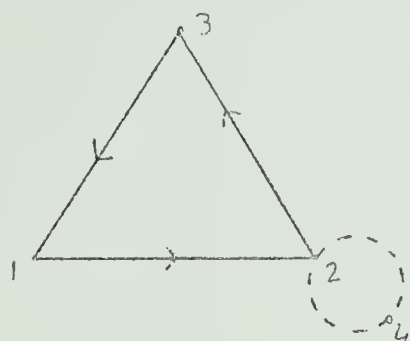
The entries are in two lines for each graph. In the first line in order the entries are: the number of vertices, the number of solid arrows (except for Table A.8 which lists the number of all arrows, solid plus two dotted), lattice constant, horizontal weight, and vertical weight or weights. The lattice constants are the f.c.c. ones.

In the second line the graph itself is listed, for example



1 2 2 3 3 1 1 2 2 1 .

Dotted arrows come first in the listing. In Tables A.6, A.7 the spurious vertex is kept in the listing of the graph, for example



2 4 4 2 2 3 3 1 1 2 .

Since $v(k) = v(\ell-k)$ we did not list all vertical weights, only 4 for $\ell \leq 7$ and 5 for $\ell = 8$.

Below is the subdivision of graphs according to the number of solid arrow ℓ .

TABLE A.1

Fluctuation-like static susceptibility graphs

ℓ	No. of graphs
1	1
2	1
3	4
4	9
5	30
6	80
7	268
8	854
Total	1247

TABLE A.2

Fluctuation-like dynamic
susceptibility graphs

ℓ	No. of graphs
3	1
4	2
5	10
6	26
7	103
8	348
Total	490

TABLE A.3

Bubble graphs

ℓ	No. of graphs
2	1
3	1
4	5
5	8
6	32
7	69
8	245
Total	361

TABLE A.4

Fourth order fluctuation graphs

ℓ	No. of graphs
2	2
3	3
4	20
5	61
6	292
7	1117
Total	1495

TABLE A.5

2	1	0.6000000000000000	01	2	2						
1	2	2	1								
3	2	0.6600000000000000	02	2	6						
1	2	2	3	3	1						
2	3	0.6000000000000000	01	2	2						
1	2	2	1	1	2	2	1				
3	3	0.6600000000000000	02	4	8						
1	2	2	1	1	3	3	1				
4	3	0.7020000000000000	03	2	24						
1	2	2	3	3	4	4	1				
4	3	-0.6000000000000000	02	4	24						
1	2	2	1	3	4	4	3				
3	4	0.8000000000000000	01	6	14						
1	2	2	1	1	2	2	3	3	1		
3	4	0.8000000000000000	01	6	8						
2	1	1	2	2	3	3	1	1	2		
3	4	0.6600000000000000	02	4	10						
2	3	3	1	1	2	2	1	1	2		
4	4	0.2400000000000000	03	4	40						
1	2	2	1	1	3	3	4	4	1		
4	4	0.7020000000000000	03	4	40						
1	3	3	4	4	1	1	2	2	1		
4	4	0.2200000000000000	03	6	40						
3	4	4	1	1	2	2	1	1	3		
5	4	0.7350000000000000	04	2	120						
1	2	2	3	3	4	4	5	5	1		
5	4	-0.2640000000000000	03	4	120						
1	2	2	1	3	4	4	5	5	2		
5	4	-0.2220000000000000	04	2	120						
3	4	4	5	5	3	1	2	2	1		
2	5	0.6000000000000000	01	2	2						
1	2	2	1	1	2	2	1	1	2	2	1
3	5	0.8000000000000000	01	6	10						
1	2	2	3	3	1	1	2	2	3	3	1
3	5	0.8000000000000000	01	6	34						
1	2	2	1	1	3	3	2	2	3	3	1
3	5	0.6600000000000000	02	4	16						
1	2	2	1	1	2	2	1	1	3	3	1
3	5	0.6600000000000000	02	4	12						
1	3	3	1	1	2	2	1	1	2	2	1
4	5	0.3300000000000000	02	3	34						
1	2	2	1	1	2	2	3	3	4	4	1
4	5	0.2300000000000000	02	3	48						
2	1	1	2	2	3	3	4	4	1	1	2
4	5	0.7020000000000000	03	4	60						
2	3	3	4	4	1	1	2	2	1	1	2
4	5	0.7020000000000000	03	2	60						
3	4	4	1	1	2	2	1	1	2	2	3
4	5	0.7020000000000000	03	4	34						
1	2	2	1	1	3	3	4	4	3	3	1

4	5	0.702000000000000	03	2	80								
1	3	3	4	4	3	3	1	1	2	2	1		
4	5	0.220000000000000	03	6	72								
1	2	2	1	1	3	3	1	1	4	4	1		
4	5	0.360000000000000	02	2	80								
1	2	2	3	3	1	1	2	2	4	4	1		
4	5	0.240000000000000	03	4	48								
2	3	3	1	1	2	2	4	4	1	1	2		
4	5	0.360000000000000	02	4	80								
1	2	2	1	1	3	3	2	2	4	4	1		
4	5	0.240000000000000	03	4	84								
1	3	3	2	2	4	4	1	1	2	2	1		
4	5	-0.600000000000000	02	4	60								
1	2	2	1	3	4	4	3	3	4	4	3		
4	5	-0.600000000000000	02	4	60								
3	4	4	3	3	4	4	3	1	2	2	1		
5	5	0.124800000000000	04	4	240								
1	2	2	1	1	3	3	4	4	5	5	1		
5	5	0.735000000000000	04	4	240								
1	3	3	4	4	5	5	1	1	2	2	1		
5	5	0.678000000000000	04	4	240								
3	4	4	5	5	1	1	2	2	1	1	3		
5	5	0.240000000000000	04	4	240								
1	2	2	3	3	1	1	4	4	5	5	1		
5	5	0.108000000000000	04	4	240								
2	3	3	1	1	4	4	5	5	1	1	2		
5	5	-0.222000000000000	04	2	240								
1	2	2	1	3	4	4	3	3	5	5	3		
5	5	-0.222000000000000	04	4	240								
3	4	4	3	3	5	5	3	1	2	2	1		
6	5	0.762660000000000	05	2	720								
1	2	2	3	3	4	4	5	5	6	6	1		
6	5	0.120200000000000	04	6	720								
1	2	2	1	3	4	4	3	5	6	6	5		
6	5	-0.141600000000000	04	4	720								
1	2	2	1	3	4	4	5	5	6	6	3		
6	5	-0.309780000000000	05	2	720								
3	4	4	5	5	6	6	3	1	2	2	1		
6	5	-0.415200000000000	04	4	720								
1	2	2	3	3	1	4	5	5	6	6	4		
3	6	0.800000000000000	01	12	24								
2	1	1	3	3	1	1	2	2	3	3	1	1	2
3	6	0.660000000000000	02	2	28								
2	3	3	1	1	2	2	1	1	3	3	1	1	2
3	6	0.800000000000000	01	6	22								
1	2	2	1	1	2	2	1	1	2	2	3	3	1
3	6	0.800000000000000	01	6	16								
2	1	1	2	2	1	1	2	2	3	3	1	1	2
3	6	0.660000000000000	02	4	14								
2	3	3	1	1	2	2	1	1	2	2	1	1	2

3	6	0.8000000000000000	01	12	44
1	2	2	1 1 3 3 1 1 2 2 3 3 1		
4	6	0.2400000000000000	03	2	206
1	2	2	1 1 2 2 3 3 4 4 3 3 1		
4	6	0.2400000000000000	03	2	114
2	1	1	2 2 3 3 4 4 3 3 1 1 2		
4	6	0.7020000000000000	03	4	140
2	3	3	4 4 3 3 1 1 2 2 1 1 2		
4	6	0.2400000000000000	03	4	154
3	4	4	3 3 1 1 2 2 1 1 2 2 3		
4	6	0.2400000000000000	03	4	176
1	2	2	1 1 3 3 1 1 2 2 4 4 1		
4	6	0.2400000000000000	03	4	96
2	1	1	3 3 1 1 2 2 4 4 1 1 2		
4	6	0.2400000000000000	03	4	112
1	3	3	1 1 2 2 4 4 1 1 2 2 1		
4	6	0.2400000000000000	03	4	112
3	1	1	2 2 4 4 1 1 2 2 1 1 3		
4	6	0.2200000000000000	03	12	112
2	4	4	1 1 2 2 1 1 3 3 1 1 2		
4	6	0.7020000000000000	03	4	112
4	1	1	2 2 1 1 3 3 1 1 2 2 4		
4	6	0.2400000000000000	03	4	112
1	2	2	1 1 2 2 1 1 3 3 4 4 1		
4	6	0.7020000000000000	03	4	84
1	3	3	4 4 1 1 2 2 1 1 2 2 1		
4	6	0.2200000000000000	03	6	84
3	4	4	1 1 2 2 1 1 2 2 1 1 3		
4	6	0.3600000000000000	02	3	114
1	2	2	3 3 1 1 2 2 3 3 4 4 1		
4	6	0.3300000000000000	02	3	63
3	1	1	2 2 3 3 4 4 1 1 2 2 3		
4	6	0.2400000000000000	03	4	70
3	4	4	1 1 2 2 3 3 1 1 2 2 3		
4	6	0.2600000000000000	02	3	206
1	2	2	1 1 3 3 2 2 3 3 4 4 1		
4	6	0.3600000000000000	02	3	212
2	1	1	3 3 2 2 3 3 4 4 1 1 2		
4	6	0.3300000000000000	02	3	204
1	3	3	2 2 3 3 4 4 1 1 2 2 1		
4	6	0.2400000000000000	03	4	233
3	4	4	1 1 2 2 1 1 3 3 2 2 3		
5	6	0.1080000000000000	04	3	504
1	2	2	1 1 3 3 1 1 4 4 5 5 1		
5	6	0.6780000000000000	04	2	504
1	4	4	5 5 1 1 2 2 1 1 3 3 1		
5	6	0.4250000000000000	03	12	504
4	5	5	1 1 2 2 1 1 3 3 1 1 4		
5	6	0.1680000000000000	03	10	588
1	2	2	1 1 2 2 3 3 4 4 5 5 1		

5	6	0.16800000000000	03	10	336								
2	1	1	2	2	3	3	4	4	5	5	1	1	2
5	6	0.73500000000000	04	4	420								
2	3	3	4	4	5	5	1	1	2	2	1	1	2
5	6	0.73500000000000	04	4	420								
3	4	4	5	5	1	1	2	2	1	1	2	2	3
5	6	0.23230000000000	04	4	560								
1	2	2	1	1	3	3	4	4	3	3	5	5	1
5	6	0.23280000000000	04	4	560								
2	1	1	3	3	4	4	3	3	5	5	1	1	2
5	6	0.73500000000000	04	2	560								
1	3	3	4	4	3	3	5	5	1	1	2	2	1
5	6	0.67800000000000	04	4	560								
3	5	5	1	1	2	2	1	1	2	3	4	4	3
5	6	0.24960000000000	04	4	588								
1	2	2	1	1	3	3	4	4	5	5	3	3	1
5	6	0.24960000000000	04	4	560								
1	3	3	4	4	5	5	3	3	1	1	2	2	1
5	6	0.73500000000000	04	4	588								
3	4	4	5	5	3	3	1	1	2	2	1	1	3
5	6	0.67800000000000	04	2	588								
4	5	5	3	3	1	1	2	2	1	1	3	3	4
5	6	0.38400000000000	03	2	560								
1	2	2	1	1	3	3	2	2	4	4	5	5	1
5	6	0.38400000000000	03	2	560								
2	1	1	3	3	2	2	4	4	5	5	1	1	2
5	6	0.12480000000000	04	4	588								
1	3	3	2	2	4	4	5	5	1	1	2	2	1
5	6	0.24960000000000	04	4	588								
2	4	4	5	5	1	1	2	2	1	1	3	3	2
5	6	0.23230000000000	04	2	588								
4	5	5	1	1	2	2	1	1	3	3	2	2	4
5	6	0.33400000000000	03	2	560								
1	2	2	3	3	1	1	2	2	4	4	5	5	1
5	6	0.12480000000000	04	4	336								
2	3	3	1	1	2	2	4	4	5	5	1	1	2
5	6	0.24960000000000	04	4	336								
2	4	4	5	5	1	1	2	2	3	3	1	1	2
5	6	0.23280000000000	04	2	336								
4	5	5	1	1	2	2	3	3	1	1	2	2	4
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3	7	0.6600000000000000	02	4	36
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3	7	0.8000000000000000	01	6	12
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4	7	0.7020000000000000	03	2	256
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4	7	0.7020000000000000	03	4	224
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4	7	0.7020000000000000	03	4	176
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4	7	0.2200000000000000	03	6	240
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4	7	0.2200000000000000	03	12	160
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4	7	0.3300000000000000	02	3	176										
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4	7	0.7020000000000000	03	4	112										
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5	7	0.3840000000000000	03	4	912										
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5	7	0.2496000000000000	04	4	560										
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5	7	0.6480000000000000	03	4	912										
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5	7	0.2323000000000000	04	2	560										
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5	7	0.3840000000000000	03	4	896										
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5	7	0.1248000000000000	04	4	768										
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5	7	0.6780000000000000	04	4	1408										
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5	7	0.6730000000000000	04	4	896										
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5	7	0.7250000000000000	04	4	896										
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5	7	0.24260000000000	04	4	396										
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5	7	0.23280000000000	04	4	912										
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5	7	0.36000000000000	02	24	1568										
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5	7	0.73500000000000	04	4	1632										
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5	7	0.19200000000000	03	4	1536										
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5	7	0.38400000000000	03	4	1568										
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5	7	0.69600000000000	03	4	1696										
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5	7	0.64800000000000	03	4	1696										
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6	7	0.21960000000000	04	3	4480
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6	7	0.16320000000000	04	8	4032										
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4	8	0.7020000000000000	03	4	552
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4	8	0.2400000000000000	03	4	496
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4	8	0.2400000000000000	03	4	476
1	3	3	1 1 2 2 4 4 2 2	3 3 1 1 2 2 1	
4	8	0.2400000000000000	03	4	936
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4	8	0.2400000000000000	03	4	684
4	1	1	2 2 1 1 2 2 1 1	3 3 2 2 3 3 4	
4	8	0.2400000000000000	03	4	684
3	4	4	1 1 2 2 1 1 2 2	1 1 3 3 2 2 3	
4	8	0.3600000000000000	02	8	576
3	2	2	3 3 4 4 1 1 2 2	1 1 2 2 1 1 3	
4	8	0.3600000000000000	02	8	572
2	3	3	4 4 1 1 2 2 1 1	2 2 1 1 3 3 2	
4	8	0.3600000000000000	02	8	920
2	1	1	2 2 1 1 3 3 2 2	3 3 4 4 1 1 2	

4	8	0.3200000000000000	02	16	568												
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4	8	0.3600000000000000	02	8	892												
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4	8	0.3600000000000000	02	3	892												
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4	8	0.3300000000000000	02	16	1064												
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4	8	0.3600000000000000	02	3	1116												
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4	8	0.3600000000000000	02	8	1072												
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4	8	0.2400000000000000	03	4	1164												
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4	8	0.3600000000000000	02	8	992												
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4	8	0.2600000000000000	02	8	1092												
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4	8	0.3300000000000000	02	16	232												
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4	8	0.3600000000000000	02	8	404												
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4	8	0.3600000000000000	02	3	176												
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4	8	0.3600000000000000	02	3	476												
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4	8	0.3600000000000000	02	8	476												
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4	8	0.24000000000000	03	4	336
2	3	3 4 4 1 1 2 2 1 1 3 3 4 4 1 1 2			
4	8	0.36000000000000	02	8	296
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4	8	0.33000000000000	02	16	332
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4	8	0.36000000000000	02	8	560
1	2	2 1 1 3 3 4 4 1 1 2 2 3 3 4 4 1			
4	8	0.24000000000000	03	4	108
4	1	1 2 2 3 3 1 1 2 2 3 3 1 1 2 2 4			
4	8	0.36000000000000	02	8	176
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4	8	0.36000000000000	02	4	236
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4	8	0.36000000000000	02	4	1900
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4	8	0.20000000000000	01	48	1304
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4	8	0.20000000000000	01	48	956
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4	8	0.36000000000000	02	4	548
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4	8	0.36000000000000	02	8	528
3	2	2 4 4 1 1 2 2 1 1 3 3 4 4 1 1 3			
4	8	0.20000000000000	01	48	512
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5	8	0.64800000000000	03	4	4448
4	2	2 4 4 5 5 1 1 2 2 1 1 3 3 1 1 4			
5	8	0.64800000000000	03	4	4896
3	1	1 4 4 2 2 4 4 5 5 1 1 2 2 1 1 3			
5	8	0.64800000000000	03	4	4352
2	4	4 5 5 1 1 2 2 1 1 3 3 1 1 4 4 2			
5	8	0.64300000000000	03	4	4464
2	1	1 3 3 1 1 4 4 2 2 4 4 5 5 1 1 2			
5	8	0.12480000000000	04	4	4400
1	4	4 2 2 4 4 5 5 1 1 2 2 1 1 3 3 1			
5	8	0.64800000000000	03	4	4896
1	3	3 1 1 4 4 2 2 4 4 5 5 1 1 2 2 1			
5	8	0.64800000000000	03	4	4432
1	2	2 1 1 3 3 1 1 4 4 2 2 4 4 5 5 1			
5	8	0.23280000000000	04	4	4824
5	4	4 1 1 2 2 1 1 3 3 1 1 4 4 2 2 5			
5	8	0.12480000000000	04	2	4432
4	2	2 5 5 4 4 1 1 2 2 1 1 3 3 1 1 4			
5	8	0.62600000000000	03	4	4352
4	1	1 2 2 1 1 3 3 1 1 4 4 2 2 5 5 4			

5	8	0.6960000000000000	03	4	4324
3	1	1 4 4 2 2 5 5 4 4	1 1 2 2 1 1	3	
5	8	0.6960000000000000	03	4	4448
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5	8	0.2496000000000000	04	4	1152
5	1	1 2 2 1 1 2 2 3 3	1 1 2 2 4 4	5	
5	8	0.1248000000000000	04	4	1152
3	1	1 2 2 4 4 5 5 1 1	2 2 1 1 2 2	3	
5	8	0.3840000000000000	03	2	864
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5	8	0.3840000000000000	03	2	2304
1	2	2 1 1 2 2 3 3 1 1	2 2 4 4 5 5	1	
5	8	0.2496000000000000	04	4	1728
5	3	3 1 1 2 2 1 1 2 2	3 3 1 1 4 4	5	
5	8	0.3840000000000000	03	4	2808
3	1	1 4 4 5 5 3 3 1 1	2 2 1 1 2 2	3	
5	8	0.1248000000000000	04	4	1584
2	3	3 1 1 4 4 5 5 3 3	1 1 2 2 1 1	2	
5	8	0.3840000000000000	03	4	1368
2	1	1 2 2 3 3 1 1 4 4	5 5 3 3 1 1	2	
5	8	0.2496000000000000	04	4	1728
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5	8	0.3840000000000000	03	4	2592
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5	8	0.2496000000000000	04	4	1584
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5	8	0.3340000000000000	03	4	2688
4	1	1 2 2 5 5 4 4 1 1	2 2 1 1 3 3	4	
5	8	0.1248000000000000	04	4	1584
3	4	4 1 1 2 2 5 5 4 4	1 1 2 2 1 1	3	
5	8	0.2328000000000000	04	4	1584
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5	8	0.3340000000000000	03	4	1328
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5	8	0.1248000000000000	04	4	1584
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5	2	2 5 5 1 1 2 2 1 1	3 3 4 4 1 1	5	
5	8	0.2040000000000000	03	16	4896
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5	8	0.2496000000000000	04	2	5472
4	1	1 5 5 2 2 5 5 1 1	2 2 1 1 3 3	4	
5	8	0.2328000000000000	04	4	1440
5	1	1 2 2 1 1 3 3 4 4	1 1 3 3 4 4	5	
5	8	0.1080000000000000	04	8	1440
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5	8	0.64800000000000	03	4	2336												
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5	8	-0.26400000000000	03	6	1152												
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5	8	-0.26400000000000	03	6	1584												
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5	8	-0.22200000000000	04	2	2016												
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5	8	-0.26400000000000	03	4	1008												
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5	8	-0.22200000000000	04	4	2520												
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5	8	-0.26400000000000	03	6	2016												
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5	8	-0.26400000000000	03	6	3523												
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5	8	-0.26400000000000	03	12	1008												
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5	8	-0.26400000000000	03	12	2016												
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5	8	-0.22200000000000	04	2	1008												
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5	8	0.49500000000000	03	24	1440												
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5	8	0.73500000000000	04	4	1008												
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5	8	0.67800000000000	04	2	2304												
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5	8	0.67800000000000	04	2	1584												
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5	8	0.73500000000000	04	2	2016												
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5	8	0.23230000000000	04	2	1584												
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5	8	0.23280000000000	04	4	3168												
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5	8	0.23280000000000	04	2	1152												
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5	8	0.23280000000000	04	4	1728												
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5	8	0.10800000000000	04	2	5472												
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5	8	0.10800000000000	04	4	864												
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5	8	0.69600000000000	03	4	4324												
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5	8	0.64800000000000	03	2	2304												
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5	8	0.24000000000000	02	6	2016												
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5	8	0.24000000000000	02	6	3803												
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5	8	0.64800000000000	03	4	2803												
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5	8	0.64800000000000	03	4	2803												
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5	8	0.64800000000000	03	4	2803												
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5	8	0.26000000000000	02	24	2688												
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5	8	0.24000000000000	02	24	2336												
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5	8	0.64800000000000	03	4	2308												
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5	8	0.24000000000000	02	24	4352												
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5	8	0.64800000000000	03	4	2880												
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5	8	0.64800000000000	03	2	2688												
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5	8	0.24000000000000	02	12	2400												
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5	8	0.36000000000000	02	12	2688												
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5	8	0.24000000000000	02	12	4448												
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5	8	0.69600000000000	03	4	2592												
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3	4	4 1 1 3 3 2 2 5 5 1 1 2 2 1 1 3			
5	8	0.2040000000000000	03	16	2432
3	2	2 5 5 1 1 2 2 1 1 3 3 4 4 1 1 3			
5	8	0.6480000000000000	03	4	2592
2	5	5 1 1 2 2 1 1 3 3 4 4 1 1 3 3 2			
5	8	0.1920000000000000	03	4	2400
2	1	1 3 3 4 4 1 1 3 3 2 2 5 5 1 1 2			
5	8	0.1920000000000000	03	4	4352
1	3	3 4 4 1 1 3 3 2 2 5 5 1 1 2 2 1			
5	8	0.1920000000000000	03	4	2336
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5	8	0.6760000000000000	03	4	5544
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5	8	0.4300000000000000	02	8	5120
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5	8	0.4800000000000000	02	3	5040
3	4	4 5 5 1 1 2 2 1 1 3 3 2 2 4 4 3			
5	8	0.3600000000000000	02	12	5112
3	2	2 4 4 3 3 4 4 5 5 1 1 2 2 1 1 3			
5	8	0.3340000000000000	03	4	5144
2	4	4 3 3 4 4 5 5 1 1 2 2 1 1 3 3 2			
5	8	0.6960000000000000	03	4	3168
5	1	1 2 2 1 1 3 3 4 4 2 2 3 3 4 4 5			
5	8	0.6960000000000000	03	4	3168
4	5	5 1 1 2 2 1 1 3 3 4 4 2 2 3 3 4			
5	8	0.3840000000000000	03	4	3072
4	2	2 3 3 4 4 5 5 1 1 2 2 1 1 3 3 4			
5	8	0.4800000000000000	02	3	5160
3	4	4 2 2 3 3 4 4 5 5 1 1 2 2 1 1 3			
5	8	0.3600000000000000	02	24	2872
2	3	3 4 4 5 5 1 1 2 2 1 1 3 3 4 4 2			
5	8	0.4800000000000000	02	3	2848
2	1	1 3 3 4 4 2 2 3 3 4 4 5 5 1 1 2			
5	8	0.3340000000000000	03	4	2864
1	3	3 4 4 2 2 3 3 4 4 5 5 1 1 2 2 1			
5	8	0.4800000000000000	02	3	2324
1	2	2 1 1 3 3 4 4 2 2 3 3 4 4 5 5 1			
5	8	0.6780000000000000	04	4	1440
5	1	1 2 2 1 1 2 2 1 1 3 3 1 1 4 4 5			
5	8	0.1080000000000000	04	3	1440
3	1	1 4 4 5 5 1 1 2 2 1 1 2 2 1 1 3			
5	8	0.1080000000000000	04	8	2160
2	1	1 2 2 1 1 3 3 1 1 4 4 5 5 1 1 2			
5	8	0.6780000000000000	04	2	2160
5	1	1 2 2 1 1 3 3 1 1 4 4 1 1 2 2 5			
5	8	0.1080000000000000	04	3	2160
4	1	1 2 2 5 5 1 1 2 2 1 1 3 3 1 1 4			

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5	8	0.42500000000000	03	24	2160
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5	8	0.10800000000000	04	8	2160
1	4	4	1 1 2 2 5 5 1 1	2 2 1 1 3 3	1
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5	8	0.73500000000000	04	4	1008
5	1	1	2 2 1 1 2 2 1 1	2 2 3 3 4 4	5
5	8	0.16300000000000	03	10	1152
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5	8	0.16800000000000	03	10	1584
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5	8	0.67300000000000	04	4	2016
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3 4 4 5 5 3 3 4 4 6 6 7 7 3 1 2 2 1
7 8 -0.57144000000000 05 8 30240
7 4 1 2 2 3 3 1 4 5 5 4 4 5 5 6 6 7
7 8 -0.25680000000000 04 16 24192
5 4 4 5 5 6 6 7 7 4 1 2 2 3 3 1 4 5
7 8 -0.25680000000000 04 16 42336
4 5 5 4 4 5 5 6 6 7 7 4 1 2 2 3 3 1
7 8 -0.21948000000000 05 8 30240
7 5 1 2 2 3 3 4 4 1 5 6 6 5 5 6 6 7
7 8 -0.25680000000000 04 12 24192
6 5 5 6 6 7 7 5 1 2 2 3 3 4 4 1 5 6
7 8 -0.25680000000000 04 12 42336
5 6 6 5 5 6 6 7 7 5 1 2 2 3 3 4 4 1

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7 8 0.1065360000000D 06 8 40320
7 1 1 2 2 1 1 3 3 4 4 5 5 3 3 6 6 7
7 8 0.2293680000000D 06 2 40320
5 3 3 6 6 7 7 1 1 2 2 1 1 3 3 4 4 5
7 8 0.2355840000000D 06 4 40320
3 6 6 7 7 1 1 2 2 1 1 3 3 4 4 5 5 3
7 8 0.2293680000000D 06 2 40320
3 4 4 5 5 3 3 6 6 7 7 1 1 2 2 1 1 3
7 8 0.3643200000000D 05 4 40320
2 1 1 3 3 4 4 5 5 3 3 6 6 7 7 1 1 2
7 8 0.2669280000000D 06 4 40320
1 3 3 4 4 5 5 3 3 6 6 7 7 1 1 2 2 1
7 8 0.3643200000000D 05 4 40320
1 2 2 1 1 3 3 4 4 5 5 3 3 6 6 7 7 1
7 8 0.1040160000000D 06 4 40320
7 1 1 2 2 1 1 3 3 4 4 5 5 6 6 3 3 7
7 8 0.4911840000000D 06 2 40320
6 3 3 7 7 1 1 2 2 1 1 3 3 4 4 5 5 6
7 8 0.5908800000000D 05 8 40320
3 7 7 1 1 2 2 1 1 3 3 4 4 5 5 6 6 3
7 8 0.4911840000000D 06 2 40320
3 4 4 5 5 6 6 3 3 7 7 1 1 2 2 1 1 3
7 8 0.3825600000000D 05 4 40320
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7 8 0.1310400000000D 06 4 40320
1 3 3 4 4 5 5 6 6 3 3 7 7 1 1 2 2 1
7 8 0.3825600000000D 05 4 40320
1 2 2 1 1 3 3 4 4 5 5 6 6 3 3 7 7 1
7 8 0.2209920000000D 06 4 40320
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7 8 0.1176720000000D 06 4 40320
6 4 4 7 7 1 1 2 2 1 1 3 3 4 4 5 5 6
7 8 0.1219200000000D 06 8 40320
4 7 7 1 1 2 2 1 1 3 3 4 4 5 5 6 6 4
7 8 0.1884000000000D 05 8 40320
2 1 1 3 3 4 4 5 5 6 6 4 4 7 7 1 1 2
7 8 0.2669280000000D 06 4 24192
7 1 1 2 2 3 3 1 1 2 2 4 4 5 5 6 6 7
7 8 0.5146800000000D 05 4 24192
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7 8 0.1646400000000D 05 2 40320
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7 8 0.2669280000000D 06 4 42336
7 4 4 1 1 2 2 3 3 1 1 4 4 5 5 6 6 7
7 8 0.2217600000000D 05 8 40320
4 1 1 2 2 3 3 1 1 4 4 5 5 6 6 7 7 4
7 8 0.1310400000000D 06 4 42336
3 1 1 4 4 5 5 6 6 7 7 4 4 1 1 2 2 3
7 8 0.1310400000000D 06 4 24192
7 1 1 2 2 3 3 4 4 1 1 2 2 5 5 6 6 7

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7 8 0.6508800000000D 05 2 40320
7 1 1 2 2 3 3 1 1 4 4 5 5 2 2 6 6 7
7 8 0.6508800000000D 05 2 40320
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7 8 0.1537200000000D 05 4 40320
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7 8 0.3132000000000D 04 4 40320
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7 8 0.6775200000000D 05 2 40320
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7 8 0.3888000000000D 04 2 40320
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7 8 0.3717600000000D 05 8 40320
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7 8 0.7900800000000D 05 4 40320
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7 8 0.2263200000000D 05 8 40320
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7 8 0.6508800000000D 05 2 40320
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7 8 0.1537200000000D 05 2 40320
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7 8 0.3132000000000D 04 2 40320
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7 8 0.6000000000000D 04 6 40320
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7 8 0.9504000000000D 04 2 40320
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8 8 -0.4792320000000D 06 4 120960
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8 8 -0.1357920000000D 06 8 120960
3 1 4 5 5 6 6 7 7 8 8 4 1 2 2 1 1 3
8 8 -0.1660656000000D 07 4 120960
8 3 1 2 2 1 3 4 4 5 5 3 3 6 6 7 7 8
8 8 -0.8089200000000D 06 4 120960
5 3 3 6 6 7 7 8 8 3 1 2 2 1 3 4 4 5
8 8 -0.7327440000000D 06 8 120960
8 4 1 2 2 3 3 1 4 5 5 4 4 6 6 7 7 8

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8 8 0.118477800000D 07      4 120960
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8 8 0.393504000000D 06      4 120960
6 5 5 7 7 8 8 5 1 2 2 1 3 4 4 3 5 6
8 8 0.314376000000D 06      8 120960
8 6 1 2 2 1 3 4 4 5 5 3 6 7 7 6 6 8
8 8 -0.119328000000D 06      8 120960
5 4 4 6 6 7 7 8 8 4 1 2 2 3 3 1 4 5
8 8 -0.299244000000D 06      8 120960
8 5 1 2 2 3 3 4 4 1 5 6 6 5 5 7 7 8
8 8 -0.992160000000D 05      8 120960
6 5 5 7 7 8 8 5 1 2 2 3 3 4 4 1 5 6
8 8 -0.451832400000D 07      4 120960
7 8 8 3 1 2 2 1 3 4 4 3 3 5 5 6 6 7
8 8 -0.442117200000D 07      2 120960
6 7 7 8 8 3 1 2 2 1 3 4 4 3 3 5 5 6
8 8 -0.479232000000D 06      4 120960
2 1 3 4 4 3 3 5 5 6 6 7 7 8 8 3 1 2
8 8 -0.623757600000D 07      2 120960
8 4 1 2 2 1 1 3 3 1 4 5 5 6 6 7 7 8
8 8 -0.137851200000D 07      4 120960
7 8 8 3 1 2 2 1 3 4 4 5 5 3 3 6 6 7
8 8 -0.333288000000D 06      4 120960
4 5 5 3 3 6 6 7 7 8 8 3 1 2 2 1 3 4
8 8 -0.127848000000D 06      8 120960
2 1 3 4 4 5 5 3 3 6 6 7 7 8 8 3 1 2
8 8 -0.670200000000D 06      8 120960
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8 8 -0.102048000000D 07      4 120960
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8 8 -0.928560000000D 05      12 120960
7 8 8 5 1 2 2 3 3 4 4 1 5 6 6 5 5 7
8 8 -0.221683200000D 07      4 120960
4 1 5 6 6 5 5 7 7 8 8 5 1 2 2 3 3 4
8 8 -0.104232000000D 06      8 120960
7 8 8 4 1 2 2 3 3 1 4 5 5 6 6 4 4 7
8 8 -0.166512000000D 06      8 120960
3 1 4 5 5 6 6 4 4 7 7 8 8 4 1 2 2 3
8 8 0.368004000000D 06      6 120960
7 8 8 5 1 2 2 1 3 4 4 3 5 6 6 5 5 7
8 8 0.393504000000D 06      8 120960
4 3 5 6 6 5 5 7 7 8 8 5 1 2 2 1 3 4
8 8 0.134496000000D 07      4 120960
5 3 6 7 7 6 6 8 8 6 1 2 2 1 3 4 4 5
8 8 0.314376000000D 06      4 120960
2 1 3 4 4 5 5 3 6 7 7 6 6 8 8 6 1 2
8 8 0.739011600000D 07      4 120960
7 8 8 1 1 2 2 1 1 3 3 4 4 5 5 6 6 7

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8 8 0.1433452800000D 08 2 120960
6 7 7 8 8 1 1 2 2 1 1 3 3 4 4 5 5 6
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8 8 0.2262720000000D 07 4 120960
7 8 8 1 1 2 2 3 3 1 1 4 4 5 5 6 6 7
8 8 0.2154528000000D 07 4 120960
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8 8 0.2043840000000D 06 4 120960
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8 8 0.1053480000000D 07 4 120960
7 8 8 1 1 2 2 3 3 4 4 1 1 5 5 6 6 7
8 8 0.5014800000000D 06 4 120960
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8 8 0.6014880000000D 06 4 120960
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8 8 0.8086074000000D 07 4 120960
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8 8 0.3756000000000D 06 4 120960
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8 8 0.2744208000000D 07 4 120960
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8 8 0.5114040000000D 06 4 120960
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8 8 0.1346208000000D 07 4 120960
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8 8 0.7843200000000D 06 4 120960
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9 8 -0.2913360000000D 06 12 362880
6 5 7 8 8 9 9 7 1 2 2 1 3 4 4 3 5 6
9 8 0.1737924000000D 08 2 362880
9 5 1 2 2 1 3 4 4 3 5 6 6 7 7 8 8 9
9 8 0.3736800000000D 06 8 362880
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9 8 0.4911408000000D 07 4 362880
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9 8 0.1870584000000D 07 4 362880
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9 8 0.2171280000000D 06 8 362880
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9 8 0.3230400000000D 06 8 362880
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9 8 -0.5924914800000D 08 2 362880
9 3 1 2 2 1 3 4 4 5 5 6 6 7 7 8 8 9
9 8 -0.4468560000000D 06 4 362880
2 1 3 4 4 5 5 6 6 7 7 8 8 9 9 3 1 2
9 8 -0.8990280000000D 07 4 362880
9 4 1 2 2 3 3 1 4 5 5 6 6 7 7 8 8 9

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9 8 -0.922896000000D 06      4 362880
3 1 4 5 5 6 6 7 7 8 8 9 9 4 1 2 2 3
9 8 -0.381223200000D 07      4 362880
9 5 1 2 2 3 3 4 4 1 5 6 6 7 7 8 8 9
9 8 -0.183835200000D 07      4 362880
4 1 5 6 6 7 7 8 8 9 9 5 1 2 2 3 3 4
9 8 0.828485220000D 08      2 362880
9 1 1 2 2 3 3 4 4 5 5 6 6 7 7 8 8 9

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TABLE A.6

A55

2	3	0.6000000000000000	01	2	1	0	0	0
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1	4	4	2 2 1 1 3 3 2 2 1					
3	4	0.8000000000000000	01	6	4	3	0	0
2	4	4	1 1 3 3 2 2 1 1 2					
2	5	0.6000000000000000	01	2	1	0	0	0
1	3	3	2 2 1 1 2 2 1 1 2 2 1					
3	5	0.6600000000000000	02	4	4	2	2	0
1	4	4	2 2 1 1 2 2 3 3 2 2 1					
3	5	0.8000000000000000	01	6	6	6	5	0
1	4	4	2 2 1 1 3 3 1 2 3 3 2					
3	5	0.8000000000000000	01	6	2	2	1	0
1	4	4	2 2 3 3 1 1 2 2 3 3 1					
4	5	-0.6900000000000000	02	4	20	8	2	0
1	5	5	2 2 1 1 2 2 1 3 4 4 3					
4	5	0.7020000000000000	03	2	14	14	12	0
3	5	5	2 2 1 1 2 2 3 3 4 4 3					
4	5	0.3300000000000000	02	8	20	16	6	0
1	5	5	2 2 1 1 2 2 3 3 4 4 1					
4	5	0.3300000000000000	02	3	10	6	8	0
2	5	5	1 1 2 2 3 3 4 4 1 1 2					
4	5	0.3600000000000000	02	4	14	14	12	0
4	5	5	2 2 3 3 4 4 1 1 2 2 4					
4	5	0.3600000000000000	02	2	16	16	8	0
3	5	5	1 1 4 4 3 3 1 1 2 2 3					
3	6	0.8000000000000000	01	6	6	3	0	4
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3	6	0.8000000000000000	01	6	4	1	2	2
2	4	4	1 1 2 2 1 1 2 2 3 3 1 1 2					
3	6	0.6600000000000000	02	2	4	4	4	4
1	4	4	2 2 3 3 2 2 3 3 1 1 3 3 1					
3	6	0.8000000000000000	01	12	4	3	4	2
1	4	4	3 3 1 1 2 2 3 3 2 2 3 3 1					
3	6	0.8000000000000000	01	12	8	7	4	6
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4	6	0.2400000000000000	03	2	44	36	16	14
2	5	5	3 3 2 2 3 3 1 1 4 4 1 1 2					
4	6	0.2400000000000000	03	2	22	14	14	14
3	5	5	2 2 3 3 1 1 4 4 1 1 2 2 3					
4	6	0.7020000000000000	03	4	22	22	18	16
3	5	5	1 1 4 4 1 1 2 2 3 3 2 2 3					
4	6	0.2400000000000000	03	4	32	28	16	24
2	5	5	3 3 4 4 2 2 3 3 2 2 1 1 2					
4	6	0.2400000000000000	03	4	16	12	16	8
3	5	5	2 2 3 3 4 4 2 2 1 1 2 2 3					
4	6	0.2400000000000000	03	4	24	14	12	12
1	5	5	4 4 1 1 4 4 1 1 2 2 3 3 1					
4	6	0.3300000000000000	02	8	34	34	26	16
1	5	5	3 3 4 4 1 1 2 2 3 3 2 2 1					

4	6	0.3600000000000000	02	8	34	32	26	28									
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4	6	0.3600000000000000	02	8	34	32	26	22									
2	5	5	3	3	4	4	1	1	3	3	2	2	1	1	2		
4	6	0.3300000000000000	02	3	10	10	10	8									
1	5	5	3	3	2	2	1	3	2	2	1	1	4	4	3		
4	6	0.3600000000000000	02	8	20	20	12	10									
2	5	5	1	1	3	3	2	2	1	1	4	4	3	3	2		
5	6	0.1680000000000000	03	10	60	40	44	48									
2	6	6	1	1	2	2	3	3	4	4	5	5	1	1	2		
5	6	0.1680000000000000	03	10	120	100	56	36									
1	6	6	2	2	3	3	4	4	5	5	1	1	2	2	1		
5	6	0.3840000000000000	03	2	96	96	64	48									
2	6	6	1	1	4	4	3	3	2	2	1	1	5	5	2		
5	6	0.3840000000000000	03	4	84	84	76	72									
1	6	6	2	2	1	1	4	4	2	3	2	2	5	5	1		
5	6	0.2496000000000000	04	4	84	84	76	72									
1	6	6	2	2	3	3	2	2	1	1	4	4	5	5	1		
5	6	0.7350000000000000	04	2	80	80	80	80									
4	6	6	1	1	2	2	1	1	3	3	4	4	5	5	4		
5	6	-0.2640000000000000	03	4	120	60	24	12									
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5	6	-0.2640000000000000	03	6	120	100	56	36									
1	6	6	2	2	1	1	2	2	3	3	1	4	5	5	4		
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5	6	0.3600000000000000	02	6	80	80	80	80									
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2	7	0.6000000000000000	01	2	1	0	0	0									
1	3	3	2	2	1	1	2	2	1	1	2	2	1	2	1		
3	7	0.8000000000000000	01	6	20	14	11	13									
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3	7	0.8000000000000000	01	12	10	10	9	9									
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3	7	0.8000000000000000	01	6	2	1	2	1									
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4	7	0.7020000000000000	03	4	32	32	26	22									
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4	7	0.7020000000000000	03	2	42	42	32	24
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4	7	0.3600000000000000	02	8	44	44	40	32
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4	7	0.3600000000000000	02	8	44	34	38	28
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4	7	0.2400000000000000	03	4	44	44	38	34
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4	7	0.3600000000000000	02	4	90	90	68	60
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4	7	0.2000000000000000	01	24	90	84	70	64
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4	7	0.3600000000000000	02	8	50	50	42	34
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4	7	0.2000000000000000	01	24	50	46	36	36
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4	7	0.2000000000000000	01	12	100	100	60	52
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4	7	-0.6900000000000000	02	4	42	12	2	0
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4	7	0.3600000000000000	02	8	48	48	30	30									
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4	7	0.2200000000000000	03	6	40	28	28	24									
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4	7	0.2400000000000000	03	2	76	72	60	60									
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5	7	0.1248000000000000	04	4	168	108	88	84									
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5	7	0.7350000000000000	04	4	204	204	192	184									
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5	7	0.6780000000000000	04	2	176	176	160	160									
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5	7	0.3840000000000000	03	4	238	226	194	190									
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5	7	0.6960000000000000	03	2	206	202	178	182									
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5	7	0.1680000000000000	03	10	238	238	198	142									
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5	7	0.3840000000000000	03	4	238	226	194	166									
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5	7	0.3600000000000000	02	24	204	204	192	184									
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5	7	0.3840000000000000	03	4	140	140	100	76									
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5	7	0.1680000000000000	03	10	70	70	70	62									
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5	7	0.3840000000000000	03	4	212	212	196	164									
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5	7	0.1920000000000000	03	4	212	204	180	172									
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5	7	0.3840000000000000	03	4	206	206	194	178									
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5	7	0.1920000000000000	03	2	206	202	182	162									
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5	7	0.1920000000000000	03	2	206	202	178	182									
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5	7	0.1920000000000000	03	2	228	228	168	128									
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5	7	0.3840000000000000	03	4	114	114	110	110									
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5	7	0.3600000000000000	02	12	136	136	104	72									
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5	7	-0.2640000000000000	03	6	84	84	64	48									
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5	7	0.6960000000000000	03	2	228	228	168	128									
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5	7	0.6480000000000000	03	2	192	192	144	144									
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5	7	0.2496000000000000	04	4	114	114	110	110									
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5	7	0.2496000000000000	04	4	206	206	194	178									
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5	7	0.1248000000000000	04	4	224	200	136	144									
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5	7	0.1248000000000000	04	4	112	88	104	80									
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5	7	0.7350000000000000	04	4	140	140	140	140									
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5	7	-0.2640000000000000	03	6	252	252	232	216									
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5	7	0.2040000000000000	03	8	308	260	152	104									
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5	7	0.2040000000000000	03	8	154	106	98	98									
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5	7	0.2496000000000000	04	4	154	154	134	118									
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5	7	0.2040000000000000	03	16	224	200	136	144									
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5	7	0.2040000000000000	03	16	112	88	104	80									
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6	7	-0.3097800000000000	05	2	588	588	548	516									
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6	7	-0.1536000000000000	04	2	672	672	512	384									
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6	7	-0.1536000000000000	04	4	588	588	548	516									
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6	7	0.1292000000000000	04	6	840	480	240	120									
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6	7	-0.1416000000000000	04	8	840	720	480	312									
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6	7	-0.1416000000000000	04	8	420	300	300	324									
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6	7	-0.2440000000000000	03	24	840	720	480	312									
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6	7	-0.2440000000000000	03	24	420	300	300	324									
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6	7	-0.1416000000000000	04	4	840	480	240	120									
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6	7	0.7626600000000000	05	2	560	560	560	560									
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6	7	0.1274400000000000	05	4	588	588	548	516									
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6	7	0.2400000000000000	04	2	672	672	512	384									
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6	7	0.2400000000000000	04	4	588	588	548	516									
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6	7	0.9660000000000000	03	4	588	588	548	516									
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6	7	0.2196000000000000	04	8	588	588	548	516									
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6	7	0.9660000000000000	03	2	672	672	512	384									
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6	7	0.9700000000000000	03	12	840	720	480	312									
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6	7	0.9700000000000000	03	12	420	300	300	324									
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6	7	0.2587200000000000	05	4	560	560	560	560									
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6	7	0.6000000000000000	03	4	560	560	560	560									
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3	8	0.6600000000000000	02	4	6	6	6	6	6
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1	4	4	3 3 1 1 2 2 3 3	1 1 2 2 1 1	2 2 1				
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3	8	0.8000000000000000	01	6	28	26	18	22	20
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4	8	0.2200000000000000	03	6	56	56	52	48	48
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4	8	0.2400000000000000	03	4	56	48	56	32	48
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4	8	0.7020000000000000	03	4	44	44	36	30	28
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4	8	0.2400000000000000	03	2	88	37	30	31	32
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4	8	0.2400000000000000	03	2	132	81	22	37	28
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4	8	0.2400000000000000	03	2	144	122	70	48	40										
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4	8	0.7020000000000000	03	4	68	68	60	54	52										
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4	8	0.2400000000000000	03	4	68	52	52	44	44										
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4	8	0.3300000000000000	02	16	76	76	60	46	52										
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4	8	0.3600000000000000	02	8	152	107	78	69	80										
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4	8	0.3600000000000000	02	8	116	112	98	92	84										
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4	8	0.3600000000000000	02	8	152	143	112	89	80										
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4	8	0.3600000000000000	02	8	152	113	108	83	80										
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4	8	0.3600000000000000	02	8	56	56	40	34	32
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4	8	0.3600000000000000	02	4	68	68	62	50	52
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4	8	0.2000000000000000	01	48	68	65	56	45	44
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5	8	0.1920000000000000	03	4	536	528	480	460	440
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5	8	0.2400000000000000	02	6	256	216	240	168	256
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5	8	0.2400000000000000	02	6	512	472	320	408	384
1	6	6	2 2 1 1 2 2 3 3	1 1 4 4 2 2 5 5 1					

5	8	0.1680000000000000	03	10	224	98	92	106	112										
2	6	6	1	1	2	2	1	1	2	2	3	3	4	4	5	5	1	1	2
5	8	0.1680000000000000	03	10	336	210	84	98	128										
1	6	6	2	2	1	1	2	2	1	1	2	2	3	3	4	4	5	5	1
5	8	0.2328000000000000	04	4	448	308	248	228	224										
2	6	6	1	1	2	2	1	1	3	3	4	4	3	3	5	5	1	1	2
5	8	0.7350000000000000	04	4	224	224	224	224	224										
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5	8	0.2328000000000000	04	4	448	308	248	228	224										
1	6	6	2	2	1	1	2	2	1	1	3	3	4	4	3	3	5	5	1
5	8	0.2496000000000000	04	4	256	256	220	188	176										
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5	8	0.2496000000000000	04	4	512	344	260	220	208										
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5	8	0.2328000000000000	04	2	256	216	240	168	256										
2	6	6	1	1	3	3	1	1	2	2	4	4	2	2	5	5	1	1	2
5	8	0.2328000000000000	04	2	512	472	320	408	384										
1	6	6	2	2	1	1	3	3	1	1	2	2	4	4	2	2	5	5	1
5	8	0.2496000000000000	04	4	352	254	204	222	240										
2	6	6	1	1	3	3	1	1	2	2	4	4	5	5	2	2	1	1	2
5	8	0.1680000000000000	03	20	448	392	276	220	208										
3	6	6	1	1	2	2	4	4	5	5	3	3	1	1	2	2	1	1	3
5	8	0.1680000000000000	03	20	224	168	172	156	144										
2	6	6	1	1	3	3	1	1	2	2	4	4	5	5	3	3	1	1	2
5	8	0.7350000000000000	04	4	312	312	296	284	280										
4	6	6	1	1	2	2	1	1	3	3	1	1	2	2	4	4	5	5	4
5	8	0.6780000000000000	04	4	312	312	296	284	280										
2	6	6	4	4	5	5	4	4	1	1	2	2	1	1	3	3	1	1	2
5	8	0.2328000000000000	04	4	312	250	268	230	216										
2	6	6	1	1	3	3	1	1	2	2	4	4	5	5	4	4	1	1	2
5	8	0.2328000000000000	04	4	624	562	404	402	368										
1	6	6	2	2	1	1	3	3	1	1	2	2	4	4	5	5	4	4	1
5	8	0.1080000000000000	04	2	368	268	240	236	240										
5	6	6	4	4	5	5	1	1	2	2	1	1	3	3	1	1	4	4	5
5	8	0.6780000000000000	04	2	368	368	320	304	304										
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5	8	0.1080000000000000	04	2	736	636	416	316	224										
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5	8	0.2496000000000000	04	4	320	320	292	276	272										
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5	8	0.2496000000000000	04	4	320	256	276	244	240										
2	6	6	1	1	2	2	3	3	1	1	4	4	5	5	4	4	1	1	2
5	8	0.2496000000000000	04	4	320	320	292	276	272										
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5	8	0.2496000000000000	04	4	640	576	428	388	368										
1	6	6	2	2	1	1	2	2	3	3	1	1	4	4	5	5	4	4	1
5	8	0.1680000000000000	03	20	280	210	200	210	216										
4	6	6	3	3	4	4	5	5	1	1	2	2	1	1	2	2	3	3	4
5	8	0.1680000000000000	03	20	560	490	360	250	208										
3	6	6	4	4	3	3	4	4	5	5	1	1	2	2	1	1	2	2	3

5	8	0.7350000000000000	04	2	280	280	280	280	280
2	6	6	3 3 4 4 3 3 4 4	5 5 1 1 2 2	1 1 2				
5	8	0.2496000000000000	04	4	424	424	408	396	392
4	6	6	3 3 1 1 2 2 1 1	2 2 3 3 4 4	5 5 4				
5	8	0.7350000000000000	04	4	424	424	388	356	344
3	6	6	1 1 2 2 1 1 2 2	3 3 4 4 5 5	4 4 3				
5	8	0.2496000000000000	04	2	424	310	280	278	280
2	6	6	1 1 2 2 3 3 4 4	5 5 4 4 3 3	1 1 2				
5	8	0.2496000000000000	04	2	348	734	500	346	288
1	6	6	2 2 1 1 2 2 3 3	4 4 5 5 4 4	3 3 1				
5	8	0.3340000000000000	03	2	352	254	204	222	240
2	6	6	1 1 2 2 1 1 3 3	2 2 4 4 5 5	1 1 2				
5	8	0.3840000000000000	03	2	352	254	204	222	240
1	6	6	2 2 1 1 2 2 1 1	3 3 2 2 4 4	5 5 1				
5	8	0.3840000000000000	03	4	352	352	328	280	256
3	6	6	1 1 2 2 3 3 4 4	5 5 1 1 2 2	1 1 3				
5	8	0.1248000000000000	04	4	352	352	316	284	272
2	6	6	3 3 4 4 5 5 1 1	2 2 1 1 3 3	1 1 2				
5	8	0.3840000000000000	03	4	352	282	296	254	224
2	6	6	1 1 3 3 1 1 2 2	3 3 4 4 5 5	1 1 2				
5	8	0.3840000000000000	03	4	352	352	316	264	240
1	6	6	3 3 1 1 2 2 3 3	4 4 5 5 1 1	2 2 1				
5	8	0.3840000000000000	03	4	704	634	444	414	432
1	6	6	2 2 1 1 3 3 1 1	2 2 3 3 4 4	5 5 1				
5	8	0.3840000000000000	03	4	320	320	292	276	272
3	6	6	1 1 2 2 4 4 3 3	5 5 1 1 2 2	1 1 3				
5	8	0.3840000000000000	03	4	320	256	276	244	240
2	6	6	1 1 3 3 1 1 2 2	4 4 3 3 5 5	1 1 2				
5	8	0.3840000000000000	03	4	320	320	292	276	272
1	6	6	3 3 1 1 2 2 4 4	3 3 5 5 1 1	2 2 1				
5	8	0.3840000000000000	03	4	640	576	428	388	368
1	6	6	2 2 1 1 3 3 1 1	2 2 4 4 3 3	5 5 1				
5	8	0.6480000000000000	03	4	544	524	464	468	448
4	6	6	2 2 4 4 5 5 1 1	2 2 1 1 3 3	1 1 4				
5	8	0.6480000000000000	03	4	544	524	472	428	416
2	6	6	4 4 5 5 1 1 2 2	1 1 3 3 1 1	4 4 2				
5	8	0.6480000000000000	03	4	544	532	472	476	416
2	6	6	1 1 3 3 1 1 4 4	2 2 4 4 5 5	1 1 2				
5	8	0.1248000000000000	04	4	544	544	488	416	416
1	6	6	4 4 2 2 4 4 5 5	1 1 2 2 1 1	3 3 1				
5	8	0.6480000000000000	03	4	544	532	480	436	448
1	6	6	2 2 1 1 3 3 1 1	4 4 2 2 4 4	5 5 1				
5	8	0.1248000000000000	04	2	536	536	492	432	440
4	6	6	2 2 5 5 4 4 1 1	2 2 1 1 3 3	1 1 4				
5	8	0.6960000000000000	03	4	536	528	476	432	408
4	6	6	1 1 2 2 1 1 3 3	1 1 4 4 2 2	5 5 4				
5	8	0.6960000000000000	03	4	536	528	480	460	440
2	6	6	1 1 3 3 1 1 4 4	2 2 5 5 4 4	1 1 2				
5	8	0.3840000000000000	03	2	128	72	112	88	64
2	6	6	1 1 2 2 3 3 1 1	2 2 4 4 5 5	1 1 2				

5	8	0.3600000000000000	02	24	312	312	296	284	280
3	6	6	1 1 4 4 3 3 5 5	1 1 2 2 1 1 2 2 3					
5	8	0.6480000000000000	03	4	312	312	296	284	280
2	6	6	3 3 1 1 4 4 3 3	5 5 1 1 2 2 1 1 2					
5	8	0.2400000000000000	02	24	312	250	268	230	216
2	6	6	1 1 2 2 3 3 1 1	4 4 3 3 5 5 1 1 2					
5	8	0.2400000000000000	02	24	624	562	404	402	368
1	6	6	2 2 1 1 2 2 3 3	1 1 4 4 3 3 5 5 1					
5	8	0.6480000000000000	03	2	320	320	304	272	256
2	6	6	3 3 5 5 1 1 2 2	1 1 3 3 4 4 1 1 2					
5	8	0.2400000000000000	02	12	320	256	272	256	192
2	6	6	1 1 3 3 4 4 1 1	2 2 3 3 5 5 1 1 2					
5	8	0.3600000000000000	02	12	320	320	304	272	256
1	6	6	3 3 4 4 1 1 2 2	3 3 5 5 1 1 2 2 1					
5	8	0.2400000000000000	02	12	640	576	416	368	448
1	6	6	2 2 1 1 3 3 4 4	1 1 2 2 3 3 5 5 1					
5	8	0.2040000000000000	03	16	288	288	272	256	224
3	6	6	2 2 5 5 1 1 2 2	1 1 3 3 4 4 1 1 3					
5	8	0.1920000000000000	03	4	288	284	256	244	256
2	6	6	1 1 3 3 4 4 1 1	3 3 2 2 5 5 1 1 2					
5	8	0.1920000000000000	03	4	576	576	432	384	416
1	6	6	3 3 4 4 1 1 3 3	2 2 5 5 1 1 2 2 1					
5	8	0.1920000000000000	03	4	288	284	256	228	224
1	6	6	2 2 1 1 3 3 4 4	1 1 3 3 2 2 5 5 1					
5	8	0.4800000000000000	02	8	616	600	544	532	536
4	6	6	3 3 4 4 5 5 1 1	2 2 1 1 3 3 2 2 4					
5	8	0.4800000000000000	02	8	616	600	552	508	488
3	6	6	4 4 5 5 1 1 2 2	1 1 3 3 2 2 4 4 3					
5	8	0.3600000000000000	02	12	616	616	568	496	520
3	6	6	2 2 4 4 3 3 4 4	5 5 1 1 2 2 1 1 3					
5	8	0.3840000000000000	03	4	616	616	576	520	488
2	6	6	4 4 3 3 4 4 5 5	1 1 2 2 1 1 3 3 2					
5	8	0.3840000000000000	03	4	352	352	340	332	320
4	6	6	2 2 3 3 4 4 5 5	1 1 2 2 1 1 3 3 4					
5	8	0.4800000000000000	02	8	704	704	556	424	384
3	6	6	4 4 2 2 3 3 4 4	5 5 1 1 2 2 1 1 3					
5	8	0.3600000000000000	02	24	352	352	320	276	272
2	6	6	3 3 4 4 5 5 1 1	2 2 1 1 3 3 4 4 2					
5	8	0.4800000000000000	02	8	352	338	300	290	288
2	6	6	1 1 3 3 4 4 2 2	3 3 4 4 5 5 1 1 2					
5	8	0.3840000000000000	03	4	352	352	324	276	256
1	6	6	3 3 4 4 2 2 3 3	4 4 5 5 1 1 2 2 1					
5	8	0.4800000000000000	02	8	352	338	304	282	272
1	6	6	2 2 1 1 3 3 4 4	2 2 3 3 4 4 5 5 1					
5	8	0.1080000000000000	04	8	320	236	224	204	192
2	6	6	1 1 2 2 1 1 3 3	1 1 4 4 5 5 1 1 2					
5	8	0.1080000000000000	04	4	240	204	240	156	240
2	6	6	1 1 3 3 1 1 4 4	1 1 2 2 5 5 1 1 2					
5	8	0.1080000000000000	04	4	480	444	336	396	288
1	6	6	2 2 1 1 3 3 1 1	4 4 1 1 2 2 5 5 1					

5	8	0.3840000000000000	03	2	384	328	128	184	256										
1	6	6	2	2	1	1	2	2	3	3	1	1	2	2	4	4	5	5	1
5	8	0.3840000000000000	03	4	384	384	276	240	240										
3	6	6	1	1	4	4	5	5	3	3	1	1	2	2	1	1	2	2	3
5	8	0.1248000000000000	04	4	192	192	180	156	144										
2	6	6	3	3	1	1	4	4	5	5	3	3	1	1	2	2	1	1	2
5	8	0.3840000000000000	03	4	192	150	156	130	112										
2	6	6	1	1	2	2	3	3	1	1	4	4	5	5	3	3	1	1	2
5	8	0.3840000000000000	03	4	384	342	240	218	224										
1	6	6	2	2	1	1	2	2	3	3	1	1	4	4	5	5	3	3	1
5	8	0.3840000000000000	03	4	352	352	288	240	224										
4	6	6	1	1	2	2	5	5	4	4	1	1	2	2	1	1	3	3	4
5	8	0.3840000000000000	03	4	176	140	152	140	112										
2	6	6	1	1	3	3	4	4	1	1	2	2	5	5	4	4	1	1	2
5	8	0.3840000000000000	03	4	352	316	232	204	224										
1	6	6	2	2	1	1	3	3	4	4	1	1	2	2	5	5	4	4	1
5	8	0.2040000000000000	03	8	608	580	504	480	480										
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5	8	0.2040000000000000	03	16	608	608	548	468	432										
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5	8	0.1248000000000000	04	4	160	160	152	144	160										
4	6	6	1	1	3	3	4	4	5	5	1	1	2	2	1	1	3	3	4
5	8	0.6480000000000000	03	4	320	320	248	200	160										
3	6	6	4	4	1	1	3	3	4	4	5	5	1	1	2	2	1	1	3
5	8	-0.2640000000000000	03	6	224	98	92	106	112										
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5	8	-0.2640000000000000	03	6	336	210	84	98	128										
3	6	6	4	4	3	3	4	4	3	3	4	4	5	5	3	1	2	2	1
5	8	-0.2640000000000000	03	12	448	406	304	282	288										
5	6	6	3	3	4	4	5	5	3	1	2	2	1	3	4	4	3	3	5
5	8	-0.2220000000000000	04	2	224	224	224	224	224										
4	6	6	5	5	3	1	2	2	1	3	4	4	3	3	5	5	3	3	4
5	8	-0.2640000000000000	03	12	224	182	200	178	160										
4	6	6	3	3	5	5	3	3	4	4	5	5	3	1	2	2	1	3	4
5	8	-0.2640000000000000	03	4	336	126	36	6	0										
5	6	6	4	4	5	5	4	4	5	5	4	1	2	2	3	3	1	4	5
5	8	-0.2640000000000000	03	6	280	210	200	210	216										
4	6	6	3	3	4	4	5	5	3	1	2	2	1	1	2	2	1	3	4
5	8	-0.2640000000000000	03	6	560	490	360	250	208										
3	6	6	4	4	3	3	4	4	5	5	3	1	2	2	1	1	2	2	1
5	8	-0.2640000000000000	03	12	560	350	200	110	80										
2	6	6	1	1	2	2	1	3	4	4	3	3	4	4	5	5	3	1	2
5	8	0.4300000000000000	02	4	616	616	592	540	504										
4	6	6	3	3	5	5	4	4	1	1	2	2	1	1	3	3	2	2	4
5	8	0.1920000000000000	03	4	616	616	560	492	472										
4	6	6	1	1	2	2	1	1	3	3	2	2	4	4	3	3	5	5	4
5	8	0.1920000000000000	03	4	616	616	552	524	504										
3	6	6	2	2	4	4	3	3	5	5	4	4	1	1	2	2	1	1	3
5	8	0.2400000000000000	02	8	616	590	516	478	472										
2	6	6	1	1	3	3	2	2	4	4	3	3	5	5	4	4	1	1	2

5	8	0.1920000000000000	03	4	624	624	580	512	480										
4	6	6	2	2	3	3	5	5	4	4	1	1	2	2	1	1	3	3	4
5	8	0.4800000000000000	02	2	624	624	568	592	592										
3	6	6	4	4	2	2	3	3	5	5	4	4	1	1	2	2	1	1	3
5	8	0.2400000000000000	02	4	624	588	508	476	480										
2	6	6	1	1	3	3	4	4	2	2	3	3	5	5	4	4	1	1	2
5	8	0.4800000000000000	02	4	336	336	316	284	272										
4	6	6	3	3	5	5	4	4	1	1	2	2	3	3	1	1	2	2	4
5	8	0.1920000000000000	03	4	336	336	316	288	272										
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5	8	0.1920000000000000	03	4	336	336	312	296	288										
3	6	6	1	1	2	2	4	4	3	3	5	5	4	4	1	1	2	2	3
5	8	0.2400000000000000	02	4	672	672	508	400	368										
1	6	6	2	2	3	3	1	1	2	2	4	4	3	3	5	5	4	4	1
5	8	0.6480000000000000	03	4	320	320	248	232	224										
1	6	6	3	3	4	4	1	1	3	3	4	4	5	5	1	1	2	2	1
5	8	0.6960000000000000	03	4	304	304	236	220	208										
4	6	6	1	1	3	3	5	5	4	4	1	1	2	2	1	1	3	3	4
5	8	0.1248000000000000	04	2	152	152	148	136	152										
3	6	6	4	4	1	1	3	3	5	5	4	4	1	1	2	2	1	1	3
5	8	0.3840000000000000	03	2	432	418	376	326	304										
4	6	6	3	3	5	5	4	4	1	1	2	2	1	1	2	2	3	3	4
5	8	0.2496000000000000	04	4	432	432	396	344	320										
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5	8	0.3840000000000000	03	2	432	418	364	370	384										
3	6	6	4	4	3	3	5	5	4	4	1	1	2	2	1	1	2	2	3
5	8	0.3840000000000000	03	2	432	314	280	274	272										
2	6	6	1	1	2	2	3	3	4	4	3	3	5	5	4	4	1	1	2
5	8	0.3840000000000000	03	2	864	746	508	330	304										
1	6	6	2	2	1	1	2	2	3	3	4	4	3	3	5	5	4	4	1
5	8	0.2496000000000000	04	4	232	232	220	216	216										
4	6	6	1	1	2	2	1	1	2	2	3	3	4	4	5	5	3	3	4
5	8	0.3840000000000000	03	2	464	464	364	284	256										
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5	8	0.3840000000000000	03	2	232	172	160	164	168										
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5	8	0.3840000000000000	03	2	464	404	284	204	160										
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5	8	0.2496000000000000	04	4	640	640	636	628	624										
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5	8	0.6960000000000000	03	4	640	614	548	522	512										
3	6	6	2	2	3	3	4	4	5	5	4	4	1	1	2	2	1	1	3
5	8	0.6960000000000000	03	4	640	614	548	486	464										
2	6	6	3	3	4	4	5	5	4	4	1	1	2	2	1	1	3	3	2
5	8	0.1248000000000000	04	2	640	640	564	448	432										
1	6	6	3	3	2	2	3	3	4	4	5	5	4	4	1	1	2	2	1
5	8	0.3840000000000000	03	4	624	606	552	494	464										
4	6	6	3	3	5	5	1	1	2	2	1	1	3	3	4	4	2	2	4
5	8	0.3840000000000000	03	2	624	612	564	516	480										
4	6	6	2	2	4	4	3	3	5	5	1	1	2	2	1	1	3	3	4

5	8	0.3840000000000000	03	4	624	606	548	526	544										
3	6	6	4	4	2	2	4	4	3	3	5	5	1	1	2	2	1	1	3
5	8	0.3840000000000000	03	2	624	612	568	532	528										
2	6	6	4	4	3	3	5	5	1	1	2	2	1	1	3	3	4	4	2
5	8	0.1680000000000000	03	10	624	624	580	500	448										
1	6	6	3	3	4	4	2	2	4	4	3	3	5	5	1	1	2	2	1
5	8	0.1248000000000000	04	2	192	192	188	168	160										
5	6	6	3	3	4	4	5	5	1	1	2	2	1	1	3	3	4	4	5
5	8	0.2496000000000000	04	4	192	192	188	180	176										
5	6	6	1	1	2	2	1	1	3	3	4	4	5	5	3	3	4	4	5
5	8	0.6960000000000000	03	4	384	384	304	236	208										
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5	8	0.2040000000000000	03	16	192	192	144	128	128										
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5	8	0.2040000000000000	03	8	192	192	144	108	96										
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5	8	0.1680000000000000	03	10	112	112	108	108	112										
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5	8	0.3840000000000000	03	4	224	224	180	136	112										
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5	8	0.3840000000000000	03	2	224	224	180	128	112										
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5	8	0.2400000000000000	02	2	288	288	144	144	288										
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5	8	0.1920000000000000	03	4	304	304	236	220	208										
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5	8	0.2040000000000000	03	8	152	152	148	136	152										
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6	8	0.7632000000000000	04	4	1344	924	744	684	672										
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6	8	0.7632000000000000	04	4	896	728	800	712	640										
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6	8	0.7632000000000000	04	4	1792	1624	1216	1128	1152										
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6	8	0.6992400000000000	05	2	1344	1344	1344	1344	1344										
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6	8	0.1178400000000000	05	4	1408	1408	1312	1280	1280										
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6	8	0.2164800000000000	05	4	1408	1408	1312	1280	1280										
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6	8	0.7626600000000000	05	4	1120	1120	1120	1120	1120										
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6	8	0.7632000000000000	04	4	1120	840	800	840	864										
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6	8	0.7632000000000000	04	4	2240	1960	1440	1000	832										
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6	8	0.7632000000000000	04	2	1120	840	800	840	864										
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6	8	0.7632000000000000	04	2	2240	1960	1440	1000	832										
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6	8	0.6850800000000000	05	2	1504	1504	1496	1480	1472										
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6	8	0.4776000000000000	05	2	1568	1568	1488	1408	1376										
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6	8	0.4776000000000000	05	2	1568	1568	1488	1408	1376										
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6	8	0.7626600000000000	05	4	1568	1568	1568	1568	1568										
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6	8	0.2587200000000000	05	4	1632	1632	1560	1496	1472										
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6	8	0.2587200000000000	05	4	1632	1632	1560	1496	1472										
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6	8	0.1056000000000000	04	2	1504	1504	1496	1480	1472										
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6	8	0.1056000000000000	04	4	1536	1536	1504	1440	1408										
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6	8	0.1056000000000000	04	2	1536	1536	1504	1440	1408										
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6	8	0.1056000000000000	04	2	1504	1504	1496	1480	1472										
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6	8	0.6528000000000000	04	2	1408	1408	1312	1280	1280										
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6	8	0.6528000000000000	04	2	1408	1408	1312	1280	1280										
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6	8	0.2040000000000000	04	8	896	728	800	712	640										
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6	8	0.2040000000000000	04	8	1792	1624	1216	1128	1152										
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6	8	0.2040000000000000	04	8	896	728	800	712	640										
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6	8	0.2040000000000000	04	8	1792	1624	1216	1128	1152										
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6	8	0.6528000000000000	04	2	1536	1536	1248	1152	1152										
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6	8	0.2587200000000000	05	4	1120	1120	1120	1120	1120										
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6	8	0.2040000000000000	04	8	1120	840	800	840	864										
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6	8	0.2040000000000000	04	8	2240	1960	1440	1000	832										
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6	8	0.1274400000000000	05	4	1232	1232	1112	992	944										
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6	8	0.2040000000000000	04	4	1232	896	800	784	784										
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6	8	0.2040000000000000	04	4	2464	2128	1432	976	832										
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6	8	0.2400000000000000	04	4	1904	1820	1616	1532	1520										
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6	8	0.2400000000000000	04	4	1904	1820	1616	1412	1328										
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6	8	0.9700000000000000	03	12	1904	1904	1664	1304	1136										
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6	8	0.2587200000000000	05	4	1568	1568	1568	1568	1568										
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6	8	0.7104000000000000	04	2	1568	1568	1488	1408	1376										
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6	8	0.7104000000000000	04	2	1568	1568	1488	1408	1376										
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6	8	0.6000000000000000	03	8	1632	1632	1560	1496	1472										
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6	8	0.6000000000000000	03	8	1632	1632	1560	1496	1472										
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6	8	0.1274400000000000	05	4	1648	1648	1576	1472	1424										
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6	8	0.3648000000000000	04	2	1648	1620	1496	1356	1296										
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6	8	0.3648000000000000	04	2	1648	1620	1472	1444	1456										
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6	8	0.7104000000000000	04	2	1792	1792	1472	1152	1024										
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6	8	0.2587200000000000	05	4	896	896	896	896	896										
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6	8	0.1274400000000000	05	4	912	912	888	880	880										
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6	8	0.3648000000000000	04	2	1824	1824	1464	1144	1024										
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6	8	0.9700000000000000	03	12	560	560	560	520	496										
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6	8	0.2400000000000000	04	4	1120	1120	890	680	608										
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6	8	0.6000000000000000	03	8	1088	1088	896	672	576										
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6	8	0.1032000000000000	04	8	1408	1408	1312	1280	1280										
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6	8	0.6000000000000000	03	6	1344	1344	1344	1344	1344										
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6	8	0.9660000000000000	03	3	1648	1648	1576	1472	1424										
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6	8	0.2016000000000000	04	2	1648	1620	1472	1444	1456										
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6	8	0.2400000000000000	04	4	1648	1648	1576	1472	1424										
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6	8	0.2016000000000000	04	2	1648	1620	1496	1356	1296										
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6	8	0.2196000000000000	04	16	1648	1648	1576	1472	1424										
4	7	7	1	1	2	2	1	1	3	3	2	2	4	4	5	5	6	6	4
6	8	0.1176000000000000	04	4	1648	1620	1496	1356	1296										
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6	8	0.1176000000000000	04	4	1648	1620	1472	1444	1456										
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6	8	0.9660000000000000	03	8	1696	1696	1600	1408	1312
2	7	7	3 3 5 5 6 6 1 1 2 2 1 1 3 3 4 4 2						
6	8	0.2016000000000000	04	2	1696	1640	1488	1400	1376
2	7	7	1 1 3 3 4 4 2 2 3 3 5 5 6 6 1 1 2						
6	8	0.2400000000000000	04	4	1696	1696	1600	1408	1312
1	7	7	3 3 4 4 2 2 3 3 5 5 6 6 1 1 2 2 1						
6	8	0.2016000000000000	04	2	1696	1640	1488	1400	1376
1	7	7	2 2 1 1 3 3 4 4 2 2 3 3 5 5 6 6 1						
6	8	0.5280000000000000	03	4	1568	1568	1488	1408	1376
2	7	7	1 1 3 3 4 4 2 2 5 5 3 3 6 6 1 1 2						
6	8	0.6000000000000000	03	8	1568	1568	1568	1568	1568
1	7	7	3 3 4 4 2 2 5 5 3 3 6 6 1 1 2 2 1						
6	8	0.5280000000000000	03	4	1568	1568	1488	1408	1376
1	7	7	2 2 1 1 3 3 4 4 2 2 5 5 3 3 6 6 1						
6	8	0.6600000000000000	04	4	1504	1504	1496	1480	1472
4	7	7	1 1 2 2 1 1 3 3 4 4 5 5 3 3 6 6 4						
6	8	0.9840000000000000	03	4	1504	1504	1496	1480	1472
3	7	7	4 4 5 5 3 3 6 6 4 4 1 1 2 2 1 1 3						
6	8	0.9840000000000000	03	2	1536	1536	1504	1440	1408
5	7	7	3 3 6 6 5 5 1 1 2 2 1 1 3 3 4 4 5						
6	8	0.6600000000000000	04	2	1536	1536	1504	1440	1408
5	7	7	1 1 2 2 1 1 3 3 4 4 5 5 3 3 6 6 5						
6	8	0.1032000000000000	04	4	1536	1536	1248	1152	1152
1	7	7	2 2 3 3 1 1 2 2 4 4 1 1 5 5 6 6 1						
6	8	0.2400000000000000	04	4	912	912	888	880	880
3	7	7	1 1 2 2 4 4 3 3 5 5 6 6 1 1 2 2 3						
6	8	0.9660000000000000	03	8	912	912	888	880	880
2	7	7	3 3 1 1 2 2 4 4 3 3 5 5 6 6 1 1 2						
6	8	0.2016000000000000	04	2	1824	1824	1464	1144	1024
1	7	7	2 2 3 3 1 1 2 2 4 4 3 3 5 5 6 6 1						
6	8	0.2196000000000000	04	16	912	912	888	880	880
4	7	7	1 1 2 2 3 3 1 1 2 2 4 4 5 5 6 6 4						
6	8	0.1176000000000000	04	4	1824	1824	1464	1144	1024
1	7	7	2 2 3 3 1 1 2 2 4 4 5 5 6 6 4 4 1						
6	8	0.6000000000000000	03	4	896	896	896	896	896
2	7	7	3 3 4 4 1 1 2 2 5 5 3 3 6 6 1 1 2						
6	8	0.5280000000000000	03	2	1792	1792	1472	1152	1024
1	7	7	2 2 3 3 4 4 1 1 2 2 5 5 3 3 6 6 1						
6	8	-0.1041600000000000	05	4	1344	924	744	684	672
4	7	7	3 3 4 4 3 3 5 5 6 6 3 1 2 2 1 3 4						
6	8	-0.1041600000000000	05	4	896	728	800	712	640
4	7	7	3 3 5 5 3 3 4 4 6 6 3 1 2 2 1 3 4						
6	8	-0.1041600000000000	05	4	1792	1624	1216	1128	1152
3	7	7	4 4 3 3 5 5 3 3 4 4 6 6 3 1 2 2 1						
6	8	-0.3097800000000000	05	4	1232	1232	1112	992	944
5	7	7	3 1 2 2 1 3 4 4 3 3 4 4 5 5 6 6 5						
6	8	-0.1041600000000000	05	2	1232	896	800	784	784
4	7	7	3 3 4 4 5 5 6 6 5 5 3 1 2 2 1 3 4						
6	8	-0.1041600000000000	05	2	2464	2128	1432	976	832
3	7	7	4 4 3 3 4 4 5 5 6 6 5 5 3 1 2 2 1						

6	8	-0.4152000000000000	04	8	1344	924	744	684	672
5	7	7	4 4 5 5 4 4 6 6	4 1 2 2 3 3	1 4 5				
6	8	-0.1041600000000000	05	4	2240	1400	800	440	320
2	7	7	1 1 2 2 1 3 4 4	3 3 5 5 6 6	3 1 2				
6	8	-0.4152000000000000	04	6	1120	840	800	840	864
5	7	7	4 4 5 5 6 6 4 1	2 2 1 1 3 3	1 4 5				
6	8	-0.4152000000000000	04	6	2240	1960	1440	1000	832
4	7	7	5 5 4 4 5 5 6 6	4 1 2 2 1 1	3 3 1				
6	8	-0.1536000000000000	04	8	1904	1820	1616	1532	1520
5	7	7	4 4 5 5 6 6 3 1	2 2 1 3 4 4	3 3 5				
6	8	-0.1536000000000000	04	8	1904	1820	1616	1412	1328
4	7	7	5 5 6 6 3 1 2 2	1 3 4 4 3 3	5 5 4				
6	8	-0.1416000000000000	04	8	1904	1904	1664	1304	1136
3	7	7	5 5 4 4 5 5 6 6	3 1 2 2 1 3	4 4 3				
6	8	-0.1416000000000000	04	8	560	560	560	520	496
5	7	7	3 3 4 4 5 5 6 6	3 1 2 2 1 3	4 4 5				
6	8	-0.1536000000000000	04	8	1120	1120	880	680	608
4	7	7	5 5 3 3 4 4 5 5	6 6 3 1 2 2	1 3 4				
6	8	-0.2440000000000000	03	24	2016	2016	1896	1776	1728
6	7	7	5 5 6 6 4 1 2 2	3 3 1 4 5 5	4 4 6				
6	8	-0.2440000000000000	03	24	672	672	552	432	384
6	7	7	4 4 5 5 6 6 4 1	2 2 3 3 1 4	5 5 6				
7	8	0.6168000000000000	04	14	3360	2520	2400	2520	2592
2	8	8	1 1 2 2 3 3 4 4	5 5 6 6 7 7	1 1 2				
7	8	0.6168000000000000	04	14	6720	5880	4320	3000	2496
1	8	8	2 2 1 1 2 2 3 3	4 4 5 5 6 6	7 7 1				
7	8	0.7868580000000000	06	2	4480	4480	4480	4480	4480
1	8	8	3 3 4 4 3 3 5 5	6 6 7 7 1 1	2 2 1				
7	8	0.7675200000000000	05	4	4704	4704	4464	4224	4128
3	8	8	1 1 2 2 1 1 3 3	4 4 5 5 6 6	7 7 3				
7	8	0.3888000000000000	04	4	4480	4480	4480	4480	4480
1	8	8	2 2 3 3 1 1 4 4	2 2 5 5 6 6	7 7 1				
7	8	0.1646400000000000	05	2	4704	4704	4464	4224	4128
2	8	8	1 1 3 3 2 2 4 4	5 5 6 6 7 7	1 1 2				
7	8	0.1646400000000000	05	2	4704	4704	4464	4224	4128
1	8	8	2 2 1 1 3 3 2 2	4 4 5 5 6 6	7 7 1				
7	8	0.1161600000000000	05	2	4704	4704	4464	4224	4128
2	8	8	1 1 3 3 4 4 2 2	5 5 6 6 7 7	1 1 2				
7	8	0.1161600000000000	05	2	4704	4704	4464	4224	4128
1	8	8	2 2 1 1 3 3 4 4	2 2 5 5 6 6	7 7 1				
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7	8	0.1310400000000000	06	4	4480	4480	4480	4480	4480
1	8	8	3 3 4 4 5 5 6 6	3 3 7 7 1 1	2 2 1				
7	8	0.1646400000000000	05	2	5376	5376	4416	3456	3072
1	8	8	2 2 3 3 1 1 2 2	4 4 5 5 6 6	7 7 1				
7	8	0.2217600000000000	05	8	4704	4704	4464	4224	4128
4	8	8	1 1 2 2 3 3 1 1	4 4 5 5 6 6	7 7 4				
7	8	0.1161600000000000	05	2	5376	5376	4416	3456	3072
1	8	8	2 2 3 3 4 4 1 1	2 2 5 5 6 6	7 7 1				

7	8	0.3132000000000000	04	4	4480	4480	4480	4480	4480
1	8	8 2 2 3 3 1 1 4 4 5 5 2 2 6 6 7 7 1							
7	8	0.3388000000000000	04	2	4480	4480	4480	4480	4480
1	8	8 2 2 3 3 1 1 4 4 5 5 6 6 2 2 7 7 1							
7	8	0.2263200000000000	05	3	4480	4480	4480	4480	4480
1	8	8 4 4 5 5 6 6 4 4 7 7 1 1 2 2 3 3 1							
7	8	0.3132000000000000	04	2	4480	4480	4480	4480	4480
1	3	8 2 2 3 3 4 4 1 1 5 5 2 2 6 6 7 7 1							
7	8	-0.8856000000000000	04	10	3360	2520	2400	2520	2592
4	8	8 3 3 4 4 5 5 6 6 7 7 3 1 2 2 1 3 4							
7	8	-0.8856000000000000	04	10	6720	5880	4320	3000	2496
3	8	8 4 4 3 3 4 4 5 5 6 6 7 7 3 1 2 2 1							
7	8	-0.4009440000000000	06	2	4480	4480	4480	4480	4480
3	8	8 5 5 6 6 5 5 7 7 3 1 2 2 1 3 4 4 3							
7	8	-0.1343040000000000	06	4	4704	4704	4464	4224	4128
5	8	8 3 1 2 2 1 3 4 4 3 3 5 5 6 6 7 7 5							
7	8	-0.5714400000000000	05	4	4704	4704	4464	4224	4128
6	8	8 4 1 2 2 3 3 1 4 5 5 4 4 6 6 7 7 6							
7	8	-0.8856000000000000	04	4	6720	4200	2400	1320	960
2	8	8 1 1 2 2 1 3 4 4 5 5 6 6 7 7 3 1 2							
7	8	-0.2016000000000000	05	2	4704	4704	4464	4224	4128
4	8	8 3 3 5 5 4 4 6 6 7 7 3 1 2 2 1 3 4							
7	8	-0.2016000000000000	05	2	4704	4704	4464	4224	4128
3	8	8 4 4 3 3 5 5 4 4 6 6 7 7 3 1 2 2 1							
7	8	-0.2016000000000000	05	2	5376	5376	4416	3456	3072
3	8	8 4 4 5 5 3 3 4 4 6 6 7 7 3 1 2 2 1							
7	8	-0.2568000000000000	04	16	3360	2520	2400	2520	2592
5	8	8 4 4 5 5 6 6 7 7 4 1 2 2 3 3 1 4 5							
7	8	-0.2568000000000000	04	16	6720	5880	4320	3000	2496
4	8	8 5 5 4 4 5 5 6 6 7 7 4 1 2 2 3 3 1							
7	8	0.8640000000000000	04	6	3360	2520	2400	2520	2592
6	8	8 5 5 6 6 7 7 5 1 2 2 1 3 4 4 3 5 6							
7	8	0.8640000000000000	04	6	6720	5880	4320	3000	2496
5	8	8 6 6 5 5 6 6 7 7 5 1 2 2 1 3 4 4 3							
7	8	0.8640000000000000	04	8	6720	4200	2400	1320	960
7	8	8 6 6 7 7 6 1 2 2 1 3 4 4 5 5 3 6 7							
7	8	-0.2568000000000000	04	12	3360	2520	2400	2520	2592
6	8	8 5 5 6 6 7 7 5 1 2 2 3 3 4 4 1 5 6							
7	8	-0.2568000000000000	04	12	6720	5880	4320	3000	2496
5	8	8 6 6 5 5 6 6 7 7 5 1 2 2 3 3 4 4 1							
7	8	-0.2784000000000000	04	4	5376	5376	4416	3456	3072
4	8	8 5 5 6 6 4 4 5 5 7 7 4 1 2 2 3 3 1							
7	8	-0.1872000000000000	04	6	4480	4480	4480	4480	4480
3	8	8 4 4 5 5 3 3 6 6 4 4 7 7 3 1 2 2 1							
7	8	-0.2784000000000000	04	8	4704	4704	4464	4224	4128
5	8	8 4 4 6 6 5 5 7 7 4 1 2 2 3 3 1 4 5							

TABLE A.7

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2	2	0.6000000000000000	01	2	1	0	0	0
1	3	3	1 1 2 2 1					
3	3	0.8000000000000000	01	6	3	1	0	0
1	4	4	1 1 2 2 3 3 1					
2	4	0.6000000000000000	01	2	1	0	0	0
1	3	3	1 1 2 2 1 1 2 2 1					
3	4	0.6600000000000000	02	1	4	0	4	0
1	4	4	1 1 2 2 1 1 3 3 1					
3	4	0.6600000000000000	02	2	4	2	2	0
2	4	4	2 1 2 2 1 1 3 3 1					
4	4	0.3300000000000000	02	8	12	6	4	0
1	5	5	1 1 2 2 3 3 4 4 1					
4	4	-0.6900000000000000	02	4	12	6	4	0
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3	5	0.8000000000000000	01	12	5	1	2	0
1	4	4	1 1 2 2 1 1 2 2 3 3 1					
3	5	0.8000000000000000	01	6	5	3	3	0
3	4	4	3 1 2 2 1 1 2 2 3 3 1					
4	5	0.2400000000000000	03	2	20	4	12	0
1	5	5	1 1 2 2 1 1 3 3 4 4 1					
4	5	0.2400000000000000	03	2	20	12	10	0
2	5	5	2 1 2 2 1 1 3 3 4 4 1					
4	5	0.2400000000000000	03	4	20	12	8	0
3	5	5	3 1 2 2 1 1 3 3 4 4 1					
5	5	-0.2640000000000000	03	4	60	36	24	0
1	6	6	1 1 2 2 1 3 4 4 5 5 3					
5	5	-0.2640000000000000	03	6	60	36	24	0
3	6	6	3 1 2 2 1 3 4 4 5 5 3					
5	5	0.1680000000000000	03	10	60	36	24	0
1	6	6	1 1 2 2 3 3 4 4 5 5 1					
2	6	0.6000000000000000	01	2	1	0	0	0
1	3	3	1 1 2 2 1 1 2 2 1 1 2 2 1					
3	6	0.6600000000000000	02	2	6	0	4	0
1	4	4	1 1 2 2 1 1 2 2 1 1 3 3 1					
3	6	0.6600000000000000	02	2	6	2	2	2
2	4	4	2 1 2 2 1 1 2 2 1 1 3 3 1					
3	6	0.6600000000000000	02	2	6	4	4	4
3	4	4	3 1 2 2 1 1 2 2 1 1 3 3 1					
3	6	0.8000000000000000	01	6	3	1	1	2
1	4	4	1 1 2 2 3 3 1 1 2 2 3 3 1					
3	6	0.8000000000000000	01	3	18	6	10	8
1	4	4	1 1 2 2 1 1 3 3 2 2 3 3 1					
4	6	0.3300000000000000	02	16	30	10	10	12
1	5	5	1 1 2 2 1 1 2 2 3 3 4 4 1					
4	6	0.3300000000000000	02	16	30	20	14	12
3	5	5	3 1 2 2 1 1 2 2 3 3 4 4 1					
4	6	0.7020000000000000	03	2	42	28	22	20
2	5	5	2 1 2 2 1 1 3 3 4 4 3 3 1					
4	6	0.2200000000000000	03	1	36	0	36	0
1	5	5	1 1 2 2 1 1 3 3 1 1 4 4 1					

4	6	0.2200000000000000	03	3	36	24	20	16									
2	5	5	2	1	2	2	1	1	3	3	1	1	4	4	1		
4	6	0.3600000000000000	02	4	24	8	8	16									
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4	6	0.3600000000000000	02	4	24	16	12	10									
3	5	5	3	1	2	2	3	3	1	1	2	2	4	4	1		
4	6	0.3600000000000000	02	4	42	14	22	20									
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4	6	0.3600000000000000	02	4	42	28	22	22									
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4	6	0.7020000000000000	03	2	42	14	22	20									
1	5	5	1	1	2	2	1	1	3	3	4	4	3	3	1		
4	6	-0.6900000000000000	02	4	30	20	14	12									
1	5	5	1	1	2	2	1	3	4	4	3	3	4	4	3		
4	6	-0.6900000000000000	02	4	30	10	2	0									
3	5	5	3	1	2	2	1	3	4	4	3	3	4	4	3		
5	6	0.1248000000000000	04	2	120	40	56	72									
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5	6	0.1248000000000000	04	2	120	80	64	60									
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5	6	0.1248000000000000	04	4	120	80	56	48									
3	6	6	3	1	2	2	1	1	3	3	4	4	5	5	1		
5	6	0.1248000000000000	04	2	120	80	56	48									
4	6	6	4	1	2	2	1	1	3	3	4	4	5	5	1		
5	6	0.2040000000000000	03	16	120	80	56	48									
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5	6	0.2040000000000000	03	4	120	40	56	72									
1	6	6	1	1	2	2	3	3	1	1	4	4	5	5	1		
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5	6	-0.2220000000000000	04	2	120	80	56	48									
1	6	6	1	1	2	2	1	3	4	4	3	3	5	5	3		
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6	6	-0.1416000000000000	04	8	360	240	168	144									
3	7	7	3	1	2	2	1	3	4	4	5	5	6	6	3		
6	6	-0.2440000000000000	03	24	360	240	168	144									
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6	6	-0.1416000000000000	04	4	360	240	168	144									
1	7	7	1	1	2	2	1	3	4	4	5	5	6	6	3		
6	6	0.1292000000000000	04	6	360	240	168	144									
1	7	7	1	1	2	2	1	3	4	4	3	5	6	6	5		
6	6	0.9700000000000000	03	12	360	240	168	144									
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3	7	0.8000000000000000	01	12	7	1	2	2									
1	4	4	1	1	2	2	1	1	2	2	1	1	2	2	3	3	1
3	7	0.8000000000000000	01	6	7	5	5	5									
3	4	4	3	1	2	2	1	1	2	2	1	1	2	2	3	3	1
3	7	0.8000000000000000	01	6	14	2	8	4									
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3	7	0.8000000000000000	01	12	14	6	7	7									
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4	7	0.2400000000000000	03	4	77	33	23	23									
1	5	5	1	1	2	2	1	1	2	2	3	3	4	4	3	3	1
4	7	0.2400000000000000	03	2	77	33	39	35									
3	5	5	3	1	2	2	1	1	2	2	3	3	4	4	3	3	1
4	7	0.2400000000000000	03	2	77	55	43	37									
4	5	5	4	1	2	2	1	1	2	2	3	3	4	4	3	3	1
4	7	0.2400000000000000	03	4	56	8	36	20									
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4	7	0.2400000000000000	03	4	56	24	24	24									
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4	7	0.2400000000000000	03	4	56	40	32	28									
4	5	5	4	1	2	2	1	1	3	3	1	1	2	2	4	4	1
4	7	0.2400000000000000	03	2	42	6	20	12									
1	5	5	1	1	2	2	1	1	2	2	1	1	3	3	4	4	1
4	7	0.2400000000000000	03	2	42	18	14	14									
2	5	5	2	1	2	2	1	1	2	2	1	1	3	3	4	4	1
4	7	0.2400000000000000	03	4	42	30	22	18									
3	5	5	3	1	2	2	1	1	2	2	1	1	3	3	4	4	1
4	7	0.3600000000000000	02	8	119	51	51	51									
1	5	5	1	1	2	2	1	1	3	3	2	2	3	3	4	4	1
4	7	0.3600000000000000	02	4	119	51	53	45									
2	5	5	2	1	2	2	1	1	3	3	2	2	3	3	4	4	1
4	7	0.3600000000000000	02	4	119	85	63	53									
4	5	5	4	1	2	2	1	1	3	3	2	2	3	3	4	4	1
4	7	0.3600000000000000	02	8	35	15	13	17									
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4	7	0.3600000000000000	02	4	35	15	9	17									
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4	7	0.3600000000000000	02	4	35	25	19	17									
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5	7	0.1080000000000000	04	2	252	36	180	108									
1	6	6	1	1	2	2	1	1	3	3	1	1	4	4	5	5	1
5	7	0.1080000000000000	04	4	252	180	148	124									
2	6	6	2	1	2	2	1	1	3	3	1	1	4	4	5	5	1
5	7	0.1080000000000000	04	4	252	180	132	108									
4	6	6	4	1	2	2	1	1	3	3	1	1	4	4	5	5	1
5	7	0.1680000000000000	03	20	210	90	70	78									
1	6	6	1	1	2	2	1	1	2	2	3	3	4	4	5	5	1
5	7	0.1680000000000000	03	20	210	150	110	90									
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5	7	0.1680000000000000	03	10	210	150	110	90									
4	6	6	4	1	2	2	1	1	2	2	3	3	4	4	5	5	1
5	7	0.2328000000000000	04	4	280	120	120	152									
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5	7	0.2328000000000000	04	4	280	200	160	144									
2	6	6	2	1	2	2	1	1	3	3	4	4	3	3	5	5	1

5	7	0.23280000000000	04	2	280	200	148	124									
5	6	6	5	1	2	2	1	1	3	3	4	4	3	3	5	5	1
5	7	0.24960000000000	04	2	294	126	138	146									
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5	7	0.24960000000000	04	2	294	210	166	146									
2	6	6	2	1	2	2	1	1	3	3	4	4	5	5	3	3	1
5	7	0.24960000000000	04	2	294	126	138	146									
3	6	6	3	1	2	2	1	1	3	3	4	4	5	5	3	3	1
5	7	0.24960000000000	04	4	294	210	154	126									
4	6	6	4	1	2	2	1	1	3	3	4	4	5	5	3	3	1
5	7	0.38400000000000	03	2	294	210	166	154									
3	6	6	3	1	2	2	1	1	3	3	2	2	4	4	5	5	1
5	7	0.38400000000000	03	4	294	210	154	126									
4	6	6	4	1	2	2	1	1	3	3	2	2	4	4	5	5	1
5	7	0.38400000000000	03	4	168	72	56	88									
1	6	6	1	1	2	2	3	3	1	1	2	2	4	4	5	5	1
5	7	0.38400000000000	03	2	168	120	92	76									
3	6	6	3	1	2	2	3	3	1	1	2	2	4	4	5	5	1
5	7	0.38400000000000	03	4	168	120	88	72									
4	6	6	4	1	2	2	3	3	1	1	2	2	4	4	5	5	1
5	7	0.24000000000000	02	12	280	120	120	152									
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5	7	0.24000000000000	02	12	280	200	148	124									
3	6	6	3	1	2	2	3	3	1	1	4	4	2	2	5	5	1
5	7	0.24000000000000	02	6	280	200	152	136									
4	6	6	4	1	2	2	3	3	1	1	4	4	2	2	5	5	1
5	7	0.38400000000000	03	4	294	126	138	146									
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5	7	-0.26400000000000	03	4	210	90	30	6									
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5	7	-0.26400000000000	03	12	210	150	110	90									
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5	7	-0.26400000000000	03	12	210	90	70	78									
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5	7	-0.26400000000000	03	6	210	150	130	126									
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6	7	0.76320000000000	04	4	840	600	440	360									
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6	7	-0.10416000000000	05	4	840	600	440	360									
1	7	7	1	1	2	2	1	3	4	4	3	3	5	5	6	6	3
6	7	-0.10416000000000	05	2	840	360	360	456									
3	7	7	3	1	2	2	1	3	4	4	3	3	5	5	6	6	3
6	7	-0.10416000000000	05	2	840	600	480	432									
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6	7	-0.10416000000000	05	4	840	600	440	360									
5	7	7	5	1	2	2	1	3	4	4	3	3	5	5	6	6	3
6	7	-0.41520000000000	04	6	840	600	440	360									
1	7	7	1	1	2	2	3	3	1	4	5	5	4	4	6	6	4

6	7	-0.4152000000000000	04	2	840	360	360	456									
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6	7	-0.4152000000000000	04	4	840	600	480	432									
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6	7	0.2040000000000000	04	4	840	360	360	456									
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6	7	0.2040000000000000	04	8	840	600	440	360									
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6	7	0.2040000000000000	04	8	840	600	440	360									
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6	7	0.2040000000000000	04	4	840	600	440	360									
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6	7	0.7632000000000000	04	2	840	360	360	456									
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6	7	0.7632000000000000	04	2	840	600	480	432									
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6	7	0.7632000000000000	04	4	840	600	440	360									
3	7	7	3	1	2	2	1	1	3	3	4	4	5	5	6	6	1
7	7	0.6168000000000000	04	14	2520	1800	1320	1080									
1	8	8	1	1	2	2	3	3	4	4	5	5	6	6	7	7	1
7	7	0.8640000000000000	04	8	2520	1800	1320	1080									
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7	7	0.8640000000000000	04	6	2520	1800	1320	1080									
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7	7	-0.8856000000000000	04	4	2520	1800	1320	1080									
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7	7	-0.8856000000000000	04	10	2520	1800	1320	1080									
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7	7	-0.2568000000000000	04	12	2520	1800	1320	1080									
1	8	8	1	1	2	2	3	3	1	4	5	5	6	6	7	7	4
7	7	-0.2568000000000000	04	16	2520	1800	1320	1080									
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2	8	0.6000000000000000	01	2	1	0	0	0	0
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3	8	0.6600000000000000	02	2	8	2	2	2	2
2	4	4	2 1 2 2 1 1 2 2	1 1 2	2 1 1 3 3 1				
3	8	0.6600000000000000	02	2	8	6	6	6	6
3	4	4	3 1 2 2 1 1 2 2	1 1 2	2 1 1 3 3 1				
3	8	0.6600000000000000	02	1	12	0	8	0	8
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3	8	0.8000000000000000	01	3	40	20	24	22	24
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3	8	0.6600000000000000	02	2	8	0	4	0	4
1	4	4	1 1 2 2 1 1 2 2	1 1 2	2 1 1 3 3 1				
3	8	0.6600000000000000	02	2	12	6	6	6	6
2	4	4	2 1 2 2 1 1 2 2	1 1 3	3 1 1 3 3 1				
3	8	0.8000000000000000	01	12	8	2	3	4	2
1	4	4	1 1 2 2 1 1 2 2	3 3 1	1 2 2 3 3 1				
3	8	0.8000000000000000	01	6	8	4	4	5	4
3	4	4	3 1 2 2 1 1 2 2	3 3 1	1 2 2 3 3 1				
3	8	0.8000000000000000	01	6	40	10	20	14	18
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4	8	0.7020000000000000	03	2	88	66	56	50	48
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4	8	0.2000000000000000	01	24	200	100	86	78	72
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4	8	0.2000000000000000	01	24	200	100	68	76	80
3	5	5	3 1 2 2 1 1 3 3	4 4 2	2 3 3 4 4 1				
4	8	0.3300000000000000	02	8	36	18	10	12	16
1	5	5	1 1 2 2 3 3 4 4	1 1 2	2 3 3 4 4 1				
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1	5	5	1 1 2 2 1 1 3 3	4 4 2	2 4 4 3 3 1				
4	8	0.3300000000000000	02	16	168	84	52	42	40
1	5	5	1 1 2 2 1 1 2 2	3 3 4	4 3 3 4 4 1				
4	8	0.7020000000000000	03	2	128	32	52	28	36
1	5	5	1 1 2 2 1 1 2 2	1 1 3	3 4 4 3 3 1				
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2	5	5	2 1 2 2 1 1 2 2	1 1 3	3 4 4 3 3 1				
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3	5	5	3 1 2 2 1 1 2 2	1 1 3	3 4 4 3 3 1				
4	8	0.7020000000000000	03	2	128	96	76	64	60
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4	8	0.2200000000000000	03	3	80	40	36	32	32										
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4	8	0.2400000000000000	03	1	304	228	188	160	152										
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4	8	0.3300000000000000	02	8	112	28	48	32	32										
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4	8	0.3300000000000000	02	16	112	56	46	42	40										
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4	8	0.3300000000000000	02	8	112	84	64	52	48										
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4	8	0.3300000000000000	02	16	56	14	14	16	16										
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4	8	0.3300000000000000	02	16	56	42	32	26	24										
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4	8	0.2000000000000000	01	24	360	180	154	136	128										
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4	8	0.3600000000000000	02	4	88	22	42	32	40										
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4	8	0.3600000000000000	02	4	88	66	56	52	48										
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4	8	0.3600000000000000	02	8	176	44	94	64	60										
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4	8	0.3600000000000000	02	8	176	88	74	74	68										
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4	8	0.3600000000000000	02	8	176	132	104	90	84
4	5	5 4 1 2 2 1 1 3 3	1 1 2 2 3 3 4 4 1						
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3	6	6 3 1 2 2 1 3 4 4	3 3 5 5 4 4 5 5 3						
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1	6	6 1 1 2 2 1 3 4 4	3 3 5 5 4 4 5 5 3						
5	8	-0.2220000000000000	04	2	336	252	228	224	224
5	6	6 5 1 2 2 1 3 4 4	3 3 4 4 3 3 5 5 3						
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4	6	6 4 1 2 2 1 3 4 4	3 3 4 4 3 3 5 5 3						
5	8	-0.2220000000000000	04	2	336	84	132	120	96
3	6	6 3 1 2 2 1 3 4 4	3 3 4 4 3 3 5 5 3						
5	8	-0.2220000000000000	04	4	336	252	192	156	144
1	6	6 1 1 2 2 1 3 4 4	3 3 4 4 3 3 5 5 3						
5	8	-0.2220000000000000	04	2	560	420	340	300	288
4	6	6 4 1 2 2 1 1 2 2	1 3 4 4 3 3 5 5 3						
5	8	-0.2220000000000000	04	1	560	280	240	280	304
3	6	6 3 1 2 2 1 1 2 2	1 3 4 4 3 3 5 5 3						
5	8	-0.2220000000000000	04	2	560	280	120	40	16
1	6	6 1 1 2 2 1 1 2 2	1 3 4 4 3 3 5 5 3						
5	8	-0.2640000000000000	03	6	168	84	60	76	88
3	6	6 3 1 2 2 1 3 4 4	5 5 3 3 4 4 5 5 3						
5	8	-0.2640000000000000	03	4	168	126	96	78	72
1	6	6 1 1 2 2 1 3 4 4	5 5 3 3 4 4 5 5 3						
5	8	0.2040000000000000	03	16	616	462	352	286	264
4	6	6 4 1 2 2 1 1 2 2	3 3 4 4 5 5 3 3 1						
5	8	0.6480000000000000	03	2	384	96	192	240	96
1	6	6 1 1 2 2 1 1 3 3	4 4 1 1 3 3 5 5 1						
5	8	0.6480000000000000	03	2	384	288	240	208	192
2	6	6 2 1 2 2 1 1 3 3	4 4 1 1 3 3 5 5 1						
5	8	0.6480000000000000	03	2	384	192	144	192	192
3	6	6 3 1 2 2 1 1 3 3	4 4 1 1 3 3 5 5 1						
5	8	0.6480000000000000	03	4	384	288	224	184	176
4	6	6 4 1 2 2 1 1 3 3	4 4 1 1 3 3 5 5 1						
5	8	0.2040000000000000	03	16	448	112	232	208	160
1	6	6 1 1 2 2 1 1 2 2	3 3 1 1 4 4 5 5 1						
5	8	0.2040000000000000	03	16	448	224	192	192	192
2	6	6 2 1 2 2 1 1 2 2	3 3 1 1 4 4 5 5 1						
5	8	0.2040000000000000	03	16	448	336	272	236	224
3	6	6 3 1 2 2 1 1 2 2	3 3 1 1 4 4 5 5 1						
5	8	0.2040000000000000	03	16	448	336	256	208	192
4	6	6 4 1 2 2 1 1 2 2	3 3 1 1 4 4 5 5 1						
5	8	0.2040000000000000	03	16	448	336	256	208	192
5	6	6 5 1 2 2 1 1 2 2	3 3 1 1 4 4 5 5 1						
5	8	0.3840000000000000	03	4	280	140	108	124	136
1	6	6 1 1 2 2 3 3 1 1	2 2 3 3 4 4 5 5 1						
5	8	0.3840000000000000	03	2	280	140	84	112	136
2	6	6 2 1 2 2 3 3 1 1	2 2 3 3 4 4 5 5 1						
5	8	0.3840000000000000	03	4	280	210	160	130	120
4	6	6 4 1 2 2 3 3 1 1	2 2 3 3 4 4 5 5 1						

5	8	0.3600000000000000	02	12	272	136	96	120	144
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2	6	6	2 1 2 2 3 3 4 4 1 1 2 2 3 3 5 5 1						
5	8	0.6960000000000000	03	2	456	228	204	220	248
1	6	6	1 1 2 2 1 1 3 3 4 4 5 5 3 3 4 4 1						
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2	6	6	2 1 2 2 1 1 3 3 4 4 5 5 3 3 4 4 1						
5	8	0.6960000000000000	03	4	456	228	164	208	224
3	6	6	3 1 2 2 1 1 3 3 4 4 5 5 3 3 4 4 1						
5	8	0.6960000000000000	03	2	456	342	272	230	216
5	6	6	5 1 2 2 1 1 3 3 4 4 5 5 3 3 4 4 1						
5	8	0.3600000000000000	02	6	816	408	400	392	400
1	6	6	1 1 2 2 1 1 3 3 4 4 2 2 5 5 3 3 1						
5	8	0.3600000000000000	02	12	816	408	372	396	416
2	6	6	2 1 2 2 1 1 3 3 4 4 2 2 5 5 3 3 1						
5	8	0.3600000000000000	02	12	816	612	476	408	392
4	6	6	4 1 2 2 1 1 3 3 4 4 2 2 5 5 3 3 1						
5	8	0.1248000000000000	04	4	560	280	200	200	208
1	6	6	1 1 2 2 1 1 2 2 3 3 4 4 3 3 5 5 1						
5	8	0.1248000000000000	04	4	560	280	200	200	208
2	6	6	2 1 2 2 1 1 2 2 3 3 4 4 3 3 5 5 1						
5	8	0.1248000000000000	04	4	560	280	240	280	304
3	6	6	3 1 2 2 1 1 2 2 3 3 4 4 3 3 5 5 1						
5	8	0.1248000000000000	04	4	560	420	340	300	288
4	6	6	4 1 2 2 1 1 2 2 3 3 4 4 3 3 5 5 1						
5	8	0.1248000000000000	04	4	560	420	324	272	256
5	6	6	5 1 2 2 1 1 2 2 3 3 4 4 3 3 5 5 1						
5	8	0.1248000000000000	04	4	448	112	232	208	160
1	6	6	1 1 2 2 1 1 3 3 1 1 2 2 4 4 5 5 1						
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2	6	6	2 1 2 2 1 1 3 3 1 1 2 2 4 4 5 5 1						
5	8	0.1248000000000000	04	4	448	336	284	252	240
3	6	6	3 1 2 2 1 1 3 3 1 1 2 2 4 4 5 5 1						
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4	6	6	4 1 2 2 1 1 3 3 1 1 2 2 4 4 5 5 1						
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5	6	6	5 1 2 2 1 1 3 3 1 1 2 2 4 4 5 5 1						
5	8	0.6480000000000000	03	2	704	176	424	320	352
1	6	6	1 1 2 2 1 1 3 3 1 1 4 4 2 2 5 5 1						
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2	6	6	2 1 2 2 1 1 3 3 1 1 4 4 2 2 5 5 1						
5	8	0.6480000000000000	03	2	704	528	436	372	352
3	6	6	3 1 2 2 1 1 3 3 1 1 4 4 2 2 5 5 1						
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4	6	6	4 1 2 2 1 1 3 3 1 1 4 4 2 2 5 5 1						
5	8	0.6960000000000000	03	4	824	412	364	376	400
1	6	6	1 1 2 2 1 1 3 3 2 2 4 4 5 5 4 4 1						
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3	6	6	3 1 2 2 1 1 3 3 2 2 4 4 5 5 4 4 1						

5	8	0.6960000000000000	03	2	824	412	372	412	392										
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5	8	0.6960000000000000	03	2	824	618	496	430	408										
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5	8	0.3840000000000000	03	4	952	476	408	408	408										
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5	8	0.3840000000000000	03	2	952	476	420	384	360										
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5	8	0.2040000000000000	03	16	616	308	204	184	184										
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5	8	0.4950000000000000	03	1	576	0	576	0	576										
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5	8	0.4950000000000000	03	4	576	432	360	288	288										
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5	8	0.6780000000000000	04	1	704	176	424	320	352										
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5	8	0.6780000000000000	04	2	704	528	436	372	352										
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5	8	0.6780000000000000	04	1	704	528	424	368	352										
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5	8	0.7350000000000000	04	2	816	408	372	396	416										
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5	8	0.7350000000000000	04	2	816	612	492	428	408										
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5	8	0.7350000000000000	04	1	816	408	400	392	400										
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5	8	0.1248000000000000	04	2	336	84	132	120	96										
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5	8	0.1248000000000000	04	2	336	168	120	112	112										
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5	8	0.1248000000000000	04	4	336	252	192	156	144										
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5	8	0.1248000000000000	04	2	336	252	192	156	144										
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5	8	0.1920000000000000	03	4	848	424	376	392	368										
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5	8	0.1920000000000000	03	2	848	424	336	392	464										
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5	8	0.1920000000000000	03	4	848	636	492	412	384										
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5	8	0.1920000000000000	03	4	824	412	364	376	400										
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5	8	0.192000000000000	03	2	824	412	372	412	392										
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5	8	0.192000000000000	03	4	824	618	480	406	384										
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5	8	0.192000000000000	03	4	456	228	164	208	224										
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5	8	0.192000000000000	03	2	456	228	204	220	248										
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5	8	0.192000000000000	03	4	456	342	268	234	224										
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6	8	-0.244000000000000	03	24	1680	1260	1080	1020	1008										
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6	8	-0.244000000000000	03	48	1680	840	600	600	624										
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6	8	-0.244000000000000	03	72	1680	1260	960	780	720										
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6	8	-0.153600000000000	04	4	1344	1008	792	656	608										
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6	8	-0.153600000000000	04	4	1344	672	480	608	704										
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6	8	-0.153600000000000	04	4	1344	1008	768	624	576										
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6	8	-0.966000000000000	04	3	2016	1512	1248	1064	992										
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6	8	-0.966000000000000	04	1	2016	504	1152	1080	864										
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6	8	-0.966000000000000	04	2	2016	1512	1152	936	864										
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6	8	-0.175800000000000	05	4	2240	1680	1360	1200	1152										
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6	8	-0.175800000000000	05	2	2240	1120	960	1120	1216										
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6	8	-0.141600000000000	04	16	1680	1260	960	780	720										
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6	8	-0.141600000000000	04	16	1680	840	600	600	624										
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6	8	-0.141600000000000	04	16	1680	1260	960	780	720										
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6	8	-0.141600000000000	04	8	1680	1260	960	780	720										
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6	8	-0.141600000000000	04	4	1680	840	360	120	48										
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6	8	-0.153600000000000	04	4	2352	1764	1416	1268	1232										
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6	8	-0.153600000000000	04	4	2352	1176	1080	1144	1168										
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6	8	-0.153600000000000	04	4	2352	1764	1344	1092	1008										
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6	8	-0.309780000000000	05	2	2352	1764	1416	1228	1168										
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6	8	-0.309780000000000	05	2	2352	1176	1080	1144	1168										
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6	8	-0.309780000000000	05	2	2352	1764	1344	1092	1008										
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6	8	0.529200000000000	04	2	2016	504	1152	1080	864										
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6	8	0.529200000000000	04	4	2016	1512	1248	1064	992										
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6	8	0.529200000000000	04	4	2016	1512	1152	936	864										
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6	8	0.529200000000000	04	2	2016	1512	1152	936	864										
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6	8	0.580800000000000	04	4	2240	1120	960	1120	1216										
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6	8	0.580800000000000	04	4	2240	1680	1360	1200	1152										
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6	8	0.580800000000000	04	4	2240	1680	1280	1040	960										
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6	8	0.127440000000000	05	2	2352	1176	1080	1144	1168										
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6	8	0.127440000000000	05	2	2352	1176	1080	1144	1168										
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6	8	0.127440000000000	05	4	2352	1764	1344	1092	1008										
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6	8	0.127440000000000	05	2	2352	1764	1344	1092	1008										
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6	8	0.114000000000000	05	4	2240	1120	960	1120	1216										
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6	8	0.114000000000000	05	4	2240	1680	1360	1200	1152										
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6	8	0.114000000000000	05	4	2240	1680	1280	1040	960										
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6	8	0.900000000000000	01	24	2240	1680	1280	1040	960										
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6	8	0.384000000000000	03	4	2240	1120	960	1120	1216										
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6	8	0.384000000000000	03	4	2240	1680	1288	1064	992										
3	7	7	3	1	2	2	3	3	1	1	4	4	5	5	2	2	6	6	1
6	8	0.384000000000000	03	4	2240	1680	1280	1040	960										
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6	8	0.384000000000000	03	8	2240	1120	960	1120	1216										
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6	8	0.384000000000000	03	4	2240	1680	1288	1064	992										
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6	8	0.384000000000000	03	4	2240	1680	1312	1136	1088										
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6	8	0.3840000000000000	03	8	2240	1680	1280	1040	960
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6	8	0.9660000000000000	03	4	1344	672	480	608	704
1	7	7 1 1 2 2 3 3 4 4	1 1 2 2 5 5	6 6 1					
6	8	0.9660000000000000	03	8	1344	1008	768	624	576
3	7	7 3 1 2 2 3 3 4 4	1 1 2 2 5 5	6 6 1					
6	8	0.7776000000000000	04	4	2240	1120	960	1120	1216
1	7	7 1 1 2 2 1 1 3 3	4 4 5 5 3 3	6 6 1					
6	8	0.7776000000000000	04	4	2240	1680	1360	1200	1152
2	7	7 2 1 2 2 1 1 3 3	4 4 5 5 3 3	6 6 1					
6	8	0.7776000000000000	04	4	2240	1120	960	1120	1216
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6	8	0.7776000000000000	04	4	2240	1680	1280	1040	960
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6	8	0.7776000000000000	04	4	2240	1680	1280	1040	960
5	7	7 5 1 2 2 1 1 3 3	4 4 5 5 3 3	6 6 1					
6	8	0.7776000000000000	04	4	2240	1680	1288	1064	992
6	7	7 6 1 2 2 1 1 3 3	4 4 5 5 3 3	6 6 1					
6	8	0.2400000000000000	04	4	1344	672	480	608	704
1	7	7 1 1 2 2 3 3 1 1	2 2 4 4 5 5	6 6 1					
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3	7	7 3 1 2 2 3 3 1 1	2 2 4 4 5 5	6 6 1					
6	8	0.2400000000000000	04	4	1344	1008	768	624	576
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5	7	7 5 1 2 2 3 3 1 1	2 2 4 4 5 5	6 6 1					
6	8	0.2196000000000000	04	8	2352	1176	1080	1144	1168
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6	8	0.2196000000000000	04	16	2352	1764	1344	1092	1008
2	7	7 2 1 2 2 3 3 1 1	4 4 5 5 6 6	4 4 1					
6	8	0.9660000000000000	03	4	2352	1176	1080	1144	1168
1	7	7 1 1 2 2 1 1 3 3	4 4 2 2 5 5	6 6 1					
6	8	0.9660000000000000	03	8	2352	1764	1344	1092	1008
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6	8	0.1632000000000000	04	4	2016	504	1152	1080	864
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2	7	7 2 1 2 2 1 1 3 3	4 4 1 1 5 5	6 6 1					
6	8	0.1632000000000000	04	16	2016	1512	1152	936	864
3	7	7 3 1 2 2 1 1 3 3	4 4 1 1 5 5	6 6 1					
6	8	0.1292000000000000	04	12	1680	1260	960	780	720
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6	8	0.12920000000000	04	6	1680	840	360	120	48										
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6	8	0.97000000000000	03	24	1680	840	600	600	624										
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6	8	0.97000000000000	03	24	1680	1260	960	780	720										
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6	8	0.97000000000000	03	24	1680	1260	960	780	720										
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7	8	0.12096000000000	05	4	6720	3360	2880	3360	3648										
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7	8	0.73614000000000	05	2	6720	5040	4080	3600	3456										
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7	8	-0.66312000000000	05	2	6720	5040	4080	3600	3456										
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7	8	-0.66312000000000	05	4	6720	5040	3840	3120	2880										
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7	8	-0.21948000000000	05	4	6720	5040	4080	3600	3456										
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7	8	-0.19032000000000	05	8	6720	5040	3840	3120	2880										
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7	8	-0.19032000000000	05	4	6720	5040	4080	3600	3456										
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7	8	-0.19032000000000	05	12	6720	5040	3840	3120	2880										
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7	8	-0.10824000000000	05	16	6720	5040	3840	3120	2880										
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7	8	-0.66312000000000	05	2	6720	5040	3840	3120	2880										
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7	8	-0.10824000000000	05	8	6720	5040	3840	3120	2880										
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7	8	0.49680000000000	04	4	6720	3360	2880	3360	3648										
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7	8	0.4968000000000000	04	16	6720	5040	3840	3120	2880
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3	8	8	3 1 2 2 1 1 3 3	4 4 5 5 6 6	7 7 1				
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4	8	8	4 1 2 2 1 1 3 3	4 4 5 5 6 6	7 7 1				
7	8	0.5146800000000000	05	2	6720	5040	3840	3120	2880
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7	8	0.7361400000000000	05	4	6720	5040	3840	3120	2880
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7	8	0.7361400000000000	05	1	6720	3360	2880	3360	3648
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8	8	0.1830000000000000	05	24	20160	15120	11520	9360	8640
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6	8	0.47760000000000	05	2	96										
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8	8	-0.62375760000000	07	4	720										
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8	8	-0.33282210000000	07	4	720										
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8	8	-0.49523460000000	07	4	720										
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8	8	0.11847780000000	07	8	720										
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8	8	0.13449600000000	07	4	720										
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5	9	0.1080000000000000	04	4	76												
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5	9	0.1080000000000000	04	8	24												
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5	9	0.1080000000000000	04	4	24												
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5	9	-0.2640000000000000	03	12	84												
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5	9	-0.2640000000000000	03	12	28												
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5	9	-0.2220000000000000	04	8	14												
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5	9	-0.2640000000000000	03	24	56												
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5	9	0.6780000000000000	04	4	24												
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5	9	0.2328000000000000	04	2	44												
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5	9	0.6780000000000000	04	4	44												
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5	9	0.7350000000000000	04	4	84												
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5	9	0.64800000000000	03	4	40
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5	9	0.10800000000000	04	8	16
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5	9	0.64800000000000	03	4	44
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5	9	0.10800000000000	04	3	24
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5	9	0.23280000000000	04	4	24
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5	9	0.38400000000000	03	2	46
2	1	3 4 4 3 3 5 5 4 4	1 1 2 1 2 2 3		
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5	9	0.12480000000000	04	4	54
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5	9	0.73500000000000	04	4	42
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5	9	0.12480000000000	04	4	34
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5	9	0.2496000000000000	04	4	36												
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5	9	0.4800000000000000	02	8	86												
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5	9	0.6960000000000000	03	4	50												
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5	9	0.6960000000000000	03	4	94												
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5	9	0.2496000000000000	04	4	36												
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5	9	0.1248000000000000	04	4	58												
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5	9	0.4950000000000000	03	48	24												
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5	9	0.4950000000000000	03	48	40												
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5	9	-0.2220000000000000	04	4	42												
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5	9	0.2328000000000000	04	4	52												
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5	9	0.2328000000000000	04	4	64												
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5	9	0.6780000000000000	04	4	52												
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5	9	0.2328000000000000	04	4	56												
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6	9	0.5808000000000000	04	8	204												
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6	9	0.5808000000000000	04	8	204												
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2	1	3 4 4 2 2 5 5 6 6	5 5 1 1 2 1 3		
6	9	0.68508000000000	05	4	206
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6	9	0.72960000000000	04	12	206
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6	9	0.68508000000000	05	4	112
4	1	5 6 6 1 1 2 2 1 1	3 3 4 1 2 2 5		
6	9	0.20304000000000	05	4	96
1	2	5 6 6 1 2 1 1 3 3	4 4 1 1 3 3 5		
6	9	0.20304000000000	05	4	96
2	1	5 6 6 1 1 2 1 3 3	4 4 1 1 3 3 5		
6	9	0.20304000000000	05	4	192
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6	9	0.19620000000000	05	12	96
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6	9	0.52920000000000	04	8	140
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6	9	0.69924000000000	05	4	140
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6	9	0.19620000000000	05	12	140
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6	9	0.21648000000000	05	4	308
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6	9	0.21648000000000	05	4	154
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6	9	0.68508000000000	05	4	154
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6	9	-0.10416000000000	05	3	154
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6	9	0.47760000000000	05	2	238
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6	9	0.76266000000000	05	4	238
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6	9	0.11400000000000	05	4	206
3	2	5	6	6	1 1 2 2 1 1 3 2 4 4 3 3 5
6	9	0.20304000000000	05	4	206
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6	9	0.47760000000000	05	2	212												
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6	9	0.72960000000000	04	12	212												
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6	9	0.20304000000000	05	4	212												
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6	9	0.11400000000000	05	4	212												
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6	9	0.25920000000000	04	4	192												
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6	9	0.23800000000000	04	24	96												
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6	9	0.20304000000000	05	4	96												
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6	9	0.12144000000000	05	2	228												
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6	9	0.11400000000000	05	4	114												
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6	9	0.47760000000000	05	2	114												
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6	9	0.20304000000000	05	4	114												
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6	9	0.7296000000000000	04	6	228												
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6	9	0.7626600000000000	05	4	70												
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6	9	0.5808000000000000	04	8	136												
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6	9	0.6992400000000000	05	4	68												
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6	9	0.2592000000000000	04	8	176												
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6	9	0.2030400000000000	05	4	176												
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6	9	0.2380000000000000	04	24	176												
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6	9	0.3024000000000000	04	4	228												
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6	9	0.2030400000000000	05	4	114												
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6	9	0.1178400000000000	05	8	114												
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6	9	0.7920000000000000	03	120	144												
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6	9	0.6992400000000000	05	4	168												
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6	9	0.69924000000000	05	4	280												
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6	9	0.69924000000000	05	4	140												
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6	9	-0.17580000000000	05	16	140												
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6	9	0.69924000000000	05	4	140												
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6	9	0.69924000000000	05	4	204												
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6	9	0.69924000000000	05	4	84												
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6	9	0.69924000000000	05	4	168												
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6	9	-0.17580000000000	05	16	168												
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6	9	0.20304000000000	05	4	156												
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6	9	0.20304000000000	05	4	192												
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6	9	0.19620000000000	05	12	156												
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6	9	0.20304000000000	05	4	176												
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6	9	0.20304000000000	05	4	152												
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6	9	0.19620000000000	05	12	176												
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6	9	0.15786000000000	05	8	176												
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6	9	0.15786000000000	05	8	152												
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6	9	0.69924000000000	05	4	140												
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6	9	0.69924000000000	05	4	224												
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6	9	0.21648000000000	05	4	140												
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6	9	0.21648000000000	05	4	224												
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6	9	0.68508000000000	05	4	140												
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6	9	0.76320000000000	04	4	126												
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6	9	0.76320000000000	04	4	168												
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6	9	0.76266000000000	05	4	126												
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6	9	-0.30978000000000	05	8	280												
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6	9	0.68508000000000	05	4	154												
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6	9	0.68508000000000	05	4	112												
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6	9	-0.30978000000000	05	8	154												
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6	9	0.76266000000000	05	4	280												
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6	9	0.76266000000000	05	4	140												
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6	9	0.76320000000000	04	4	140												
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6	9	0.72960000000000	04	12	166												
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6	9	0.72960000000000	04	6	228												
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6	9	0.68508000000000	05	4	178												
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6	9	0.72960000000000	04	6	206												
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6	9	0.47760000000000	05	2	180												
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6	9	0.47760000000000	05	2	176												
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6	9	0.47760000000000	05	2	212												
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6	9	0.47760000000000	05	2	232												
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6	9	0.69924000000000	05	4	180												
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6	9	0.47760000000000	05	2	204												
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6	9	0.47760000000000	05	2	206												
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6	9	0.47760000000000	05	2	188												
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6	9	0.47760000000000	05	2	178												
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6	9	0.69924000000000	05	4	204												
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6	9	0.76266000000000	05	4	280												
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6	9	0.30240000000000	04	8	172												
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6	9	0.12144000000000	05	2	206												
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6	9	0.71040000000000	04	2	180
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6	9	0.12144000000000	05	2	172
2	5	4 6 6 1 1 2 2 3 3	1 1 4 4 2 5 4		
6	9	0.30240000000000	04	4	206
2	5	6 1 1 2 2 3 3 1 1	4 4 2 5 4 4 6		
6	9	0.71040000000000	04	2	173
1	2	4 3 3 5 5 2 2 6 6	1 2 3 3 1 1 4		
6	9	0.12144000000000	05	2	166
1	4	3 5 5 2 2 6 6 1 1	2 2 3 3 1 4 3		
6	9	-0.96600000000000	04	12	84
3	5	4 6 6 3 1 2 2 1 3	4 4 3 5 3 3 4		
6	9	0.76266000000000	05	4	140
1	3	5 4 4 3 3 6 6 1 1	2 2 1 3 4 4 5		
6	9	0.68508000000000	05	4	204
3	4	6 1 1 2 2 1 1 3 4	5 5 4 4 3 3 6		
6	9	0.69924000000000	05	4	140
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6	9	0.69924000000000	05	4	280
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6	9	0.68508000000000	05	4	188
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6	9	0.68508000000000	05	4	178
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6	9	0.25872000000000	05	4	154
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6	9	0.76266000000000	05	4	168
1	2	4 5 5 6 6 4 4 3 3	1 2 1 1 3 3 4		
6	9	0.76266000000000	05	4	238
1	2	6 4 4 3 3 1 2 1 1	3 3 4 4 5 5 6		
6	9	0.25872000000000	05	4	280
1	2	4 3 3 1 2 1 1 3 3	4 4 5 5 6 6 4		
6	9	0.76266000000000	05	4	280
1	3	4 5 5 6 6 4 4 3 3	1 1 2 2 1 3 4		
6	9	0.76266000000000	05	4	154
1	3	6 4 4 3 3 1 1 2 2	1 3 4 4 5 5 6		
6	9	0.71040000000000	04	2	178
1	2	5 3 3 6 6 1 2 3 3	1 1 4 4 2 2 5		
6	9	0.12144000000000	05	2	166
1	4	2 5 5 3 3 6 6 1 1	2 2 3 3 1 4 2		
6	9	0.30240000000000	04	4	206
1	4	5 3 3 6 6 1 1 2 2	3 3 1 4 2 2 5		
6	9	0.71040000000000	04	2	208
1	2	5 4 4 6 6 1 2 3 3	1 1 4 4 2 2 5		
6	9	0.71040000000000	04	2	180
2	3	1 4 4 2 2 5 5 4 4	6 6 1 1 2 3 1		
6	9	0.12144000000000	05	2	212
2	3	5 4 4 6 6 1 1 2 3	1 1 4 4 2 2 5		

6	9	-0.9660000000000000	04	12	168												
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6	9	-0.1041600000000000	05	4	336												
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6	9	-0.1041600000000000	05	4	168												
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6	9	-0.1041600000000000	05	8	280												
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6	9	-0.3097800000000000	05	8	280												
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6	9	-0.4152000000000000	04	24	280												
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6	9	-0.4152000000000000	04	24	140												
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6	9	-0.1758000000000000	05	16	140												
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6	9	-0.1758000000000000	05	16	140												
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6	9	-0.1041600000000000	05	4	252												
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6	9	-0.1041600000000000	05	4	168												
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6	9	0.5292000000000000	04	8	224												
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6	9	0.2164800000000000	05	4	140												
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6	9	0.5292000000000000	04	8	140												
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6	9	0.2164800000000000	05	4	224												
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6	9	0.6850800000000000	05	4	140												
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6	9	0.6992400000000000	05	4	140												
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6	9	-0.1041600000000000	05	8	224												
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6	9	0.6992400000000000	05	4	84												
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6	9	0.6850800000000000	05	4	84												
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6	9	0.5292000000000000	04	8	84												
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6	9	0.52920000000000	04	8	112												
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6	9	0.21648000000000	05	4	112												
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6	9	0.25872000000000	05	4	280												
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6	9	0.12744000000000	05	4	126												
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6	9	0.12744000000000	05	4	168												
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6	9	0.25872000000000	05	4	126												
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6	9	0.68508000000000	05	4	154												
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6	9	0.68508000000000	05	4	112												
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6	9	0.25872000000000	05	4	336												
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6	9	0.25872000000000	05	4	336												
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6	9	0.25872000000000	05	4	252												
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6	9	-0.41520000000000	04	12	252												
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6	9	0.71040000000000	04	2	204												
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6	9	0.71040000000000	04	2	180												
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6	9	0.47760000000000	05	2	204												
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6	9	0.71040000000000	04	2	178												
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6	9	0.71040000000000	04	2	232												
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6	9	0.47760000000000	05	2	178												
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6	9	0.76266000000000	05	4	140												
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6	9	0.11400000000000	05	4	188												
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6	9	0.11400000000000	05	4	212												
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6	9	0.68508000000000	05	4	188												
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6	9	0.58080000000000	04	8	204												
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6	9	0.58080000000000	04	8	180												
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6	9	0.11784000000000	05	8	206												
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7 9 -0.4009440000000D 06 4 420
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7 9 0.7361400000000D 05 8 420
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7 9 -0.20160000000000 05 4 588
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7 9 -0.66312000000000 05 8 588
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7 9 -0.12496800000000 06 2 588
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7 9 -0.12496800000000 06 2 588
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7 9 -0.36787200000000 06 4 588
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7 9 -0.20160000000000 05 4 672
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7 9 -0.66312000000000 05 8 336
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7 9 -0.13430400000000 06 8 336
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7 9 -0.12496800000000 06 4 336
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7 9 -0.12496800000000 06 2 672
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7 9 -0.36787200000000 06 4 336
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7 9 -0.48553200000000 06 8 420
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7 9 -0.57144000000000 05 4 420
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7 9 0.86400000000000 04 24 420
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7 9 -0.57144000000000 05 12 840
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7 9 -0.57144000000000 05 12 420
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9 9 0.49114080000000 07 8 5040
4 5 6 7 7 8 8 9 9 4 1 2 2 3 3 1 5 6
9 9 0.27825120000000 07 8 5040
4 5 7 8 8 9 9 4 1 2 2 3 3 1 5 6 6 7
9 9 0.37368000000000 06 8 5040
1 2 3 4 4 1 5 6 6 7 7 8 8 9 9 5 2 3
9 9 0.18705840000000 07 8 5040
5 6 7 8 8 9 9 5 1 2 2 3 3 4 4 1 6 7
9 9 -0.59249148000000 08 4 5040
1 2 3 4 4 5 5 6 6 7 7 8 8 9 9 1 2 3
9 9 -0.76616940000000 08 4 5040
1 2 4 5 5 6 6 7 7 8 8 9 9 1 2 3 3 4

```


Table A9

2	2	0.600000000000000D	01	1	2
1	2	2	1		
3	3	0.800000000000000D	01	2	6
1	2	2	3	3	1
2	4	0.600000000000000D	01	1	2
1	2	2	1	1	2
3	4	0.660000000000000D	02	1	8
1	2	2	1	1	3
4	4	0.330000000000000D	02	2	24
1	2	2	3	3	4
4	4	-0.690000000000000D	02	1	24
1	2	2	1	3	4
4	5	0.240000000000000D	03	2	40
1	2	2	1	1	3
3	5	0.800000000000000D	01	6	10
1	2	2	1	1	2
5	5	0.168000000000000D	03	2	120
1	2	2	3	3	4
5	5	-0.264000000000000D	03	2	120
1	2	2	1	3	4
2	6	0.600000000000000D	01	1	2
1	2	2	1	1	2
3	6	0.660000000000000D	02	2	12
1	2	2	1	1	2
3	6	0.800000000000000D	01	2	6
1	2	2	3	3	1
3	6	0.800000000000000D	01	1	36
1	2	2	1	1	3
4	6	0.702000000000000D	03	1	84
1	2	2	1	1	3
4	6	0.220000000000000D	03	1	72
1	2	2	1	1	3
4	6	0.360000000000000D	02	2	48
1	2	2	3	3	1
4	6	-0.690000000000000D	02	2	60
1	2	2	1	3	4
4	6	0.360000000000000D	02	2	84
1	2	2	1	1	3
4	6	0.330000000000000D	02	8	60
1	2	2	1	1	2
5	6	0.204000000000000D	03	4	240
1	2	2	3	3	1
5	6	-0.222000000000000D	04	1	240
1	2	2	1	3	4
5	6	0.124800000000000D	04	2	240
1	2	2	1	1	3
6	6	0.970000000000000D	03	2	720
1	2	2	3	3	4
6	6	-0.244000000000000D	03	4	720
1	2	2	3	3	1

6	6	-0.14160000000000D	04	2	720
1	2	2	1 3 4 4 5 5 6 6 3		
6	6	0.12920000000000D	04	1	720
1	2	2	1 3 4 4 3 5 6 6 5		
3	7	0.80000000000000D	01	6	14
1	2	2	1 1 2 2 1 1 2 2 3 3 1		
3	7	0.80000000000000D	01	6	28
1	2	2	1 1 3 3 1 1 2 2 3 3 1		
4	7	0.24000000000000D	03	2	154
1	2	2	1 1 2 2 3 3 4 4 3 3 1		
4	7	0.24000000000000D	03	4	112
1	2	2	1 1 3 3 1 1 2 2 4 4 1		
4	7	0.24000000000000D	03	2	84
1	2	2	1 1 2 2 1 1 3 3 4 4 1		
4	7	0.36000000000000D	02	4	70
1	2	2	3 3 1 1 2 2 3 3 4 4 1		
4	7	0.36000000000000D	02	4	238
1	2	2	1 1 3 3 2 2 3 3 4 4 1		
5	7	0.24000000000000D	02	6	560
1	2	2	3 3 1 1 4 4 2 2 5 5 1		
5	7	0.10800000000000D	04	2	504
1	2	2	1 1 3 3 1 1 4 4 5 5 1		
5	7	0.16800000000000D	03	10	420
1	2	2	1 1 2 2 3 3 4 4 5 5 1		
5	7	0.23280000000000D	04	2	560
1	2	2	1 1 3 3 4 4 3 3 5 5 1		
5	7	0.24960000000000D	04	2	588
1	2	2	1 1 3 3 4 4 5 5 3 3 1		
5	7	0.38400000000000D	03	2	336
1	2	2	3 3 1 1 2 2 4 4 5 5 1		
5	7	0.38400000000000D	03	2	588
1	2	2	1 1 3 3 2 2 4 4 5 5 1		
5	7	-0.26400000000000D	03	2	420
1	2	2	3 3 1 4 5 5 4 4 5 5 4		
5	7	-0.26400000000000D	03	6	420
1	2	2	1 3 4 4 3 3 4 4 5 5 3		
6	7	0.20400000000000D	04	4	1680
1	2	2	3 3 1 1 4 4 5 5 6 6 1		
6	7	0.76320000000000D	04	2	1680
1	2	2	1 1 3 3 4 4 5 5 6 6 1		
6	7	-0.10416000000000D	05	2	1680
1	2	2	1 3 4 4 3 3 5 5 6 6 3		
6	7	-0.41520000000000D	04	2	1680
1	2	2	3 3 1 4 5 5 4 4 6 6 4		
7	7	0.61680000000000D	04	2	5040
1	2	2	3 3 4 4 5 5 6 6 7 7 1		
7	7	-0.25680000000000D	04	4	5040
1	2	2	3 3 1 4 5 5 6 6 7 7 4		
7	7	-0.88560000000000D	04	2	5040
1	2	2	1 3 4 4 5 5 6 6 7 7 3		

7	7	0.86400000000000D	04	2	5040
1	2	2 1 3 4 4 3 5 6 6 7 7 5			
2	8	0.60000000000000D	01	1	2
1	2	2 1 1 2 2 1 1 2 2 1 1 2 2 1			
3	8	0.80000000000000D	01	3	80
1	2	2 1 1 2 2 1 1 3 3 2 2 3 3 1			
3	8	0.66000000000000D	02	2	16
1	2	2 1 1 2 2 1 1 2 2 1 1 3 3 1			
3	8	0.66000000000000D	02	1	24
1	2	2 1 1 2 2 1 1 3 3 1 1 3 3 1			
3	8	0.80000000000000D	01	6	16
1	2	2 1 1 2 2 3 3 1 1 2 2 3 3 1			
4	8	-0.69000000000000D	02	1	280
1	2	2 1 1 2 2 1 3 4 4 3 3 4 4 3			
4	8	0.70200000000000D	03	1	176
1	2	2 1 1 3 3 1 1 2 2 4 4 2 2 1			
4	8	-0.69000000000000D	02	2	112
1	2	2 1 3 4 4 3 3 4 4 3 3 4 4 3			
4	8	0.20000000000000D	01	12	400
1	2	2 1 1 3 3 4 4 2 2 3 3 4 4 1			
4	8	0.33000000000000D	02	2	72
1	2	2 3 3 4 4 1 1 2 2 3 3 4 4 1			
4	8	0.33000000000000D	02	1	800
1	2	2 1 1 3 3 4 4 2 2 4 4 3 3 1			
4	8	0.33000000000000D	02	4	336
1	2	2 1 1 2 2 3 3 4 4 3 3 4 4 1			
4	8	0.70200000000000D	03	2	256
1	2	2 1 1 2 2 1 1 3 3 4 4 3 3 1			
4	8	0.22000000000000D	03	3	160
1	2	2 1 1 2 2 1 1 3 3 1 1 4 4 1			
4	8	0.36000000000000D	02	8	192
1	2	2 1 1 2 2 3 3 1 1 4 4 3 3 1			
4	8	0.36000000000000D	02	2	128
1	2	2 1 1 2 2 3 3 1 1 2 2 4 4 1			
4	8	0.24000000000000D	03	2	96
1	2	2 1 1 3 3 4 4 1 1 3 3 4 4 1			
4	8	0.24000000000000D	03	1	608
1	2	2 1 1 3 3 1 1 4 4 2 2 4 4 1			
4	8	0.33000000000000D	02	8	224
1	2	2 1 1 3 3 1 1 2 2 4 4 3 3 1			
4	8	0.33000000000000D	02	8	112
1	2	2 1 1 2 2 1 1 2 2 3 3 4 4 1			
4	8	0.20000000000000D	01	6	720
1	2	2 1 1 3 3 2 2 4 4 3 3 4 4 1			
4	8	0.36000000000000D	02	2	176
1	2	2 1 1 2 2 1 1 3 3 2 2 4 4 1			
4	8	0.36000000000000D	02	8	352
1	2	2 1 1 3 3 1 1 2 2 3 3 4 4 1			
5	8	0.69600000000000D	03	2	912
1	2	2 1 1 3 3 4 4 5 5 3 3 4 4 1			

5	8	0.36000000000000D	02	6	1632
1	2	2	1	1	3 3 4 4 2 2 5 5 3 3 1
5	8	0.12480000000000D	04	4	1120
1	2	2	1	1	2 2 3 3 4 4 3 3 5 5 1
5	8	0.12480000000000D	04	4	896
1	2	2	1	1	3 3 1 1 2 2 4 4 5 5 1
5	8	-0.26400000000000D	03	2	336
1	2	2	1	3	4 4 5 5 3 3 4 4 5 5 3
5	8	-0.22200000000000D	04	1	1120
1	2	2	1	1	2 2 1 3 4 4 3 3 5 5 3
5	8	-0.22200000000000D	04	2	672
1	2	2	1	3	4 4 3 3 4 4 3 3 5 5 3
5	8	0.64800000000000D	03	2	1408
1	2	2	1	1	3 3 1 1 4 4 2 2 5 5 1
5	8	-0.26400000000000D	03	1	2016
1	2	2	1	3	4 4 3 3 5 5 4 4 5 5 3
5	8	0.69600000000000D	03	2	1648
1	2	2	1	1	3 3 2 2 4 4 5 5 4 4 1
5	8	0.38400000000000D	03	2	1904
1	2	2	1	1	3 3 2 2 3 3 4 4 5 5 1
5	8	0.20400000000000D	03	8	1232
1	2	2	1	1	2 2 3 3 4 4 5 5 3 3 1
5	8	0.64800000000000D	03	2	768
1	2	2	1	1	3 3 4 4 1 1 3 3 5 5 1
5	8	0.20400000000000D	03	16	896
1	2	2	1	1	2 2 3 3 1 1 4 4 5 5 1
5	8	0.38400000000000D	03	2	560
1	2	2	3	3	1 1 2 2 3 3 4 4 5 5 1
5	8	0.36000000000000D	02	6	544
1	2	2	3	3	4 4 1 1 2 2 3 3 5 5 1
5	8	0.49500000000000D	03	1	1152
1	2	2	1	1	3 3 1 1 4 4 1 1 5 5 1
5	8	0.67800000000000D	04	1	1408
1	2	2	1	1	3 3 1 1 4 4 5 5 4 4 1
5	8	0.73500000000000D	04	1	1632
1	2	2	1	1	3 3 4 4 5 5 4 4 3 3 1
5	8	0.12480000000000D	04	2	672
1	2	2	1	1	2 2 1 1 3 3 4 4 5 5 1
5	8	0.19200000000000D	03	2	1696
1	2	2	1	1	3 3 4 4 2 2 3 3 5 5 1
5	8	0.19200000000000D	03	2	1648
1	2	2	1	1	3 3 2 2 4 4 3 3 5 5 1
5	8	0.19200000000000D	03	2	912
1	2	2	3	3	1 1 2 2 4 4 3 3 5 5 1
6	8	0.38400000000000D	03	4	4480
1	2	2	3	3	1 1 4 4 2 2 5 5 6 6 1
6	8	0.38400000000000D	03	2	4480
1	2	2	3	3	1 1 4 4 5 5 2 2 6 6 1
6	8	0.90000000000000D	01	6	4480
1	2	2	3	3	4 4 1 1 5 5 3 3 6 6 1

6	8	0.52920000000000D	04	2	4032
1	2	2	1 1 3 3 1 1 4 4 5 5 6 6 1		
6	8	0.58080000000000D	04	2	4480
1	2	2	1 1 3 3 4 4 5 5 4 4 6 6 1		
6	8	0.12744000000000D	05	2	4704
1	2	2	1 1 3 3 4 4 5 5 6 6 3 3 1		
6	8	0.11400000000000D	05	2	4480
1	2	2	1 1 3 3 4 4 3 3 5 5 6 6 1		
6	8	-0.30978000000000D	05	1	4704
1	2	2	1 3 4 4 3 3 5 5 6 6 5 5 3		
6	8	-0.15360000000000D	04	2	4704
1	2	2	1 3 4 4 3 3 5 5 4 4 6 6 3		
6	8	-0.14160000000000D	04	2	3360
1	2	2	1 1 2 2 1 3 4 4 5 5 6 6 3		
6	8	-0.14160000000000D	04	8	3360
1	2	2	1 3 4 4 3 3 4 4 5 5 6 6 3		
6	8	-0.17580000000000D	05	1	4480
1	2	2	1 1 3 3 1 4 5 5 4 4 6 6 4		
6	8	0.96600000000000D	03	2	2688
1	2	2	3 3 4 4 1 1 2 2 5 5 6 6 1		
6	8	-0.96600000000000D	04	1	4032
1	2	2	1 3 4 4 3 3 5 5 3 3 6 6 3		
6	8	0.77760000000000D	04	4	4480
1	2	2	1 1 3 3 4 4 5 5 3 3 6 6 1		
6	8	0.24000000000000D	04	2	2688
1	2	2	3 3 1 1 2 2 4 4 5 5 6 6 1		
6	8	0.21960000000000D	04	4	4704
1	2	2	3 3 1 1 4 4 5 5 6 6 4 4 1		
6	8	0.96600000000000D	03	2	4704
1	2	2	1 1 3 3 4 4 2 2 5 5 6 6 1		
6	8	0.24000000000000D	04	2	4704
1	2	2	1 1 3 3 2 2 4 4 5 5 6 6 1		
6	8	-0.15360000000000D	04	2	2688
1	2	2	1 3 4 4 5 5 3 3 4 4 6 6 3		
6	8	-0.24400000000000D	03	24	3360
1	2	2	3 3 1 4 5 5 4 4 5 5 6 6 4		
6	8	0.16320000000000D	04	4	4032
1	2	2	1 1 3 3 4 4 1 1 5 5 6 6 1		
6	8	0.12920000000000D	04	3	3360
1	2	2	1 3 4 4 3 5 6 6 5 5 6 6 5		
6	8	0.97000000000000D	03	12	3360
1	2	2	1 1 2 2 3 3 4 4 5 5 6 6 1		
7	8	0.12096000000000D	05	4	13440
1	2	2	3 3 1 1 4 4 5 5 6 6 7 7 1		
7	8	0.49680000000000D	04	4	13440
1	2	2	3 3 4 4 1 1 5 5 6 6 7 7 1		
7	8	-0.10824000000000D	05	4	13440
1	2	2	1 3 4 4 5 5 3 3 6 6 7 7 3		
7	8	0.51468000000000D	05	2	13440
1	2	2	1 1 3 3 4 4 5 5 6 6 7 7 1		

7	8	-0.1903200000000D	05	4	13440										
1	2	2	3	3	1	4	5	5	4	4	6	6	7	7	4
7	8	-0.2194800000000D	05	2	13440										
1	2	2	1	1	3	3	1	4	5	5	6	6	7	7	4
7	8	-0.6631200000000D	05	2	13440										
1	2	2	1	3	4	4	3	3	5	5	6	6	7	7	3
7	8	0.7361400000000D	05	1	13440										
1	2	2	1	3	4	4	3	5	6	6	5	5	7	7	5
8	8	0.5304000000000D	05	2	40320										
1	2	2	1	3	4	4	3	5	6	6	7	7	8	8	5
8	8	-0.1584000000000D	05	4	40320										
1	2	2	3	3	1	4	5	5	6	6	7	7	8	8	4
8	8	0.4206900000000D	05	2	40320										
1	2	2	3	3	4	4	5	5	6	6	7	7	8	8	1
8	8	-0.6065400000000D	05	2	40320										
1	2	2	1	3	4	4	5	5	6	6	7	7	8	8	3
8	8	-0.3003150000000D	05	1	40320										
1	2	2	1	3	4	4	3	5	6	6	5	7	8	8	7
8	8	0.1830000000000D	05	4	40320										
1	2	2	1	3	4	4	5	5	3	6	7	7	8	8	6
8	8	-0.6667500000000D	04	4	40320										
1	2	2	3	3	4	4	1	5	6	6	7	7	8	8	5

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